Abstract

With the advent of new and proposed experiment facilities in the U.S. (BNL/AGS and RHIC) and elsewhere (CERN/SPS and LHC), the relativistic nuclear physics has become one of the top national priorities in nuclear physics research. We propose a new research program toward the understanding of the phenomenology of ultrarelativistic heavy-ion collisions. Our focus is to investigate the quark-gluon substructure of nuclear matter and the possibility of quark-gluon plasma formation. We have planned an extended and rigorous research program to analyze and understand experimental data, to find a promising and reliable set of signatures for the phase transition, and predict new results from fundamental theories. The following research directions and topics have been suggested to carry out the program: Heavy quark production at RHIC and LHC, Nuclear gluon structure function at small x, Disoriented chiral condensate, $U(1)_A$ particle production, Electromagnetic probe (dilepton production), QCD-based Effective field theory of multiparticle production, Wavelet analysis of multiparticle correlations, and Rapidity gaps.

Contents

Chapter 1

Progress Report

The University of Arizona research in nuclear theory performed under grant DE-FG02-93ER40792 related to relativistic heavy ion physics involve a combination of basic theory and new phenomenological ideas applied to collisions involving relativistic nuclei with hadrons or other nuclei. In the past three years we have completed following projects:

1.1 Charm Production at RHIC & LHC to $O(\alpha_s^3)$

(I. Sarcevic and P. Valerio)

The main goal of the future heavy-ion colliders, such as RHIC and LHC, is to study the properties of nuclear matter under extreme conditions, and in particular to search for the formation of a new state of matter, the quark-gluon plasma [1]. With the assumption that high energy nuclear collisions lead to the thermalized system, RHIC and LHC energies are sufficiently large to produce very dense matter, of the order of several GeV$/fm^3$, well above the critical density necessary to create the quark-gluon plasma [2]. The problem of finding a clean, detectable signal for this phase is currently one of the most challenging theoretical problems. In the last few years, the possible signals, such as thermal photons, dileptons and J/Ψ suppression, have been investigated in detail and found to have yields comparable in magnitude with those expected from a simple extrapolation of hadronic collisions [3]. Recently, open charm production in heavy-ion collisions has been proposed as an elegant method for probing the possible formation of the quark-gluon plasma [4]. However, in order to determine whether the enchanced charm production at RHIC can be unambiguously interpreted as a signal of QGP, one needs to have control of the other sources of charm production. The standard competing process is charm production through the hard collisions of partons inside the nuclei. In addition to being relevant as a background for the signal of quark-gluon plasma, this type of charm production is useful tool for studying the perturbative aspects of strong interactions and for determining the nuclear screening/shadowing effect on the gluon distribution in a nucleus.

1.1.1 Charm Quark Pair Production to $O(\alpha_s^3)$

Recently we have completed the calculation of the rapidity and transverse momentum distributions of inclusive charm quark production in p-p and Au-Au collisions at RHIC and LHC, including the $O(\alpha_s^3)$ radiative corrections and the nuclear shadowing effect [5].

In perturbative QCD differential inclusive distribution of charm production in nuclear collisions is obtained by convolution of parton densities in nuclei with a hard scattering parton cross section. For the parton cross sections, we have included leading-order subprocesses, $O(\alpha_s^2)$, such as $q+\bar{q} \rightarrow Q+\bar{Q}$ and $g+g \rightarrow Q+\bar{Q}$, and next-to-leading order contributions, $O(\alpha_s^3)$, such as $q+\bar{q} \rightarrow Q+\bar{Q}+g$, $g+q \rightarrow Q+\bar{Q}+g$, $g+\bar{q} \rightarrow Q+\bar{Q}+\bar{q}$ and $g+g \rightarrow Q+\bar{Q}+g$. The double differential inclusive distribution of charm production in central AA collisions can be written as

$$\frac{dN_c}{d^2p_T dy} = T_{AA}(0)\frac{d\sigma_c}{d^2p_T dy}, \tag{1.1}$$

where the double differential inclusive cross section is given by

$$\frac{d\sigma_c}{d^2p_T dy} = \sum_{i,j}^{partons} \int dx_a dx_b F_i^A(x_a, Q^2) F_j^A(x_b, Q^2) \frac{d\hat{\sigma}_{i,j}(Q^2, m_c, \hat{s})}{d^2p_T dy}, \tag{1.2}$$

and $T_{AA}(0)$ is the nuclear overlapping density at zero impact parameter, $F_i^A(x, Q^2)$ is the parton structure function in a nucleus, x_a and x_b are the fractional momenta of the incoming partons, $\hat{s}$ is the parton-parton c.m. energy ($\hat{s} = x_a x_b s$). The parton differential cross section calculated to $O(\alpha_s^3)$ and the next-to-leading order expression for the coupling constant $\alpha_s(Q^2)$ can be found in Ref. 6.

To obtain the number of nucleon-nucleon collisions per unit of transverse area at fixed impact parameter, we consider the nuclear overlapping function [7]

$$T_{AA}(b) = \int d^2b_1 T_A(\mid \vec{b}_1 \mid) T_A(\mid \vec{b} - \vec{b}_1 \mid), \tag{1.3}$$

where $T_A(b)$ is the nuclear profile function (i.e. the nuclear density integrated over the longitudinal size). For the nuclear density we take the Woods-Saxon distribution [8]. For central collisions the overlapping function can be approximated by $T_{AA}(0) = A^2/\pi R_A^2$, which gives $T_{Au-Au}(0) = 30.7 mb^{-1}$.

1.1.2 The Nuclear Shadowing Effect

The nuclear parton distribution, if nucleons were independent, would be given as A times the parton structure function in a nucleon. However, at high energies, the parton densities become so large that the sea quarks and gluons overlap spatially and the nucleus can not be viewed as a collection of uncorrelated nucleons. This happens when the longitudinal size of the parton, in the infinite momentum frame of the nucleus, becomes larger than the size of the nucleon. Partons from different nucleons start to interact and through annihilation effectively reduce the parton density in a nucleus. When partons inside the nucleus completely overlap, there reach a saturation point. Motivated by this simple parton picture of the nuclear shadowing effect and taking into account the $A^{1/3}$ dependence obtained by considering the modified, nonlinear Altarelli-Parisi equations with gluon recombination included, the modifying factor to the parton structure function in a nucleus can be written as [9]

$$R(x,A) \equiv \frac{F_i^A(x,Q^2)}{AF_i^N(x,Q^2)} = \begin{cases} 1-\frac{3}{16}x+\frac{3}{80} & .2 < x \leq 1 \\ 1 & x_n < x \leq .2 \\ 1-D(A^{1/3}-1)\frac{1/x-1/x_n}{1/x_A-1/x_n} & x_A \leq x \leq x_n \\ 1-D(A^{1/3}-1) & 0 < x < x_A \end{cases} \tag{1.4}$$

where $F_i^N(x,Q^2)$ is the parton structure function in a nucleon, $x_n = 1/(2r_p m_p)$, x_A is a saturation point ($x_A = 1/(2R_A m_p)$), m_p is the mass of the proton, r_p is the radius of a proton and R_A is the radius of the nucleus. By fitting the parameter D to all deep inelastic lepton-nucleus data on the ratio $F_2^A(x,Q^2)/F_2^D(x,Q^2)$ [10] we find that $R(x,A)$ has much steeper x-dependence than the data, especially for $0.002 \leq x \leq 0.1$, the region of relevance to charm production at RHIC and LHC energies. In addition, the onset of saturation in the ratio of structure functions for Xe to D is observed at much smaller value of x than predicted by the parton recombination model [11]. Thus, it is not surprising that even the best fit of Eq. (1.4) overestimates the observed shadowing effect by about 15%. Consequently charm production calculated with this shadowing function would be underestimated by about 40%.

We calculate charm production using the shadowing function that has recently been proposed as the best fit to EMC, NMC and E665 data [10] and is given by [11]

$$R(x,A) = \begin{cases} \alpha_3 - \alpha_4 x & x_0 < x \leq 0.6 \\ (\alpha_3 - \alpha_4 x_0)\frac{1+k_q\alpha_2(1/x-1/x_0)}{1+k_q A^{\alpha_1}(1/x-1/x_0)} & x \leq x_0 \end{cases} \tag{1.5}$$

The values for the parameters k_q, α_1, α_2, α_3 and x_0 can be found in Ref. 11.

Presently there is no theory which can quantitatively describe the observed nuclear shadowing effect. In Section 2.2 we outline the first theoretical calculation of the gluon density in a nucleus based on perturbative QCD.

By integrating Eq. (1.2) and multipling by $T_{AA}(0)$ (for central collisions) or by A^2/σ_{in}^{AA} (for inelastic collisions), we obtain rapidity distribution, transverse momentum distribution and the total cross section for the inclusive charm quark production in AA collisions. Comparison of our results for the total charm cross section with the low-energy hadronic and nuclear data ($20GeV \leq \sqrt{s} \leq 55GeV$) [12] indicate very good agreement for the renormalization scale $Q = m_c$ [13]. When $Q = 2m_c$ is used, our cross sections are slightly smaller than the measured values.

In our calculation we have used the two-loop evolved parton structure functions, MRS-S0 with $\Lambda_5 = 140MeV$ [13]. By using two other sets of structure functions, MRS-D0 and MRS-D– [13], we find that theoretical uncertainty due to the choice of the nucleon structure function is only 10% at RHIC. This is not surprising because the average x-value probed with charm production at $\sqrt{s} = 200GeV$ is about 10^{-2}, which is still within the range of x for which there is deep inelastic lepton-nucleon scattering data. However, at LHC energies, where charm production is probing very small values of x ($\langle x \rangle \sim 10^{-4}$), the results are very sensitive to the gluon structure function in the low-x region [14], far below the range that has currently been explored by deep-inelastic scattering experiments. In particular, we find that the total cross section for charm production is about a factor of six larger when we use MRS-D– structure function, which has "singular" behavior at low-x (i.e. $G(x) \sim x^{-1.5}$), instead of the MRS-S0 set [14]. In case of the rapidity distribution, the results for charm production obtained using MRS-D- structure function are about factor of 4.6 larger for $|y| \leq 1$. In addition, at the LHC the theoretical uncertainty due to the choice of scale is about 50%.

Our results for the rapidity distribution of inclusive charm production in central Au-Au collisions at RHIC and LHC are presented in Fig. 1 in Ref. 5 (solid line). We also show the rapidity distribution when nuclear shadowing is not included (dotted line) and the leading-order results with shadowing (short-dashed line) and without shadowing (long-dashed line). We note the the *shape* of the rapidity distribution does not seem to be affected by the next-to-leading corrections or by the nuclear shadowing effect. We find that, in the central rapidity region, the number of charm quark pairs produced per unit rapidity in central Au-Au collisions at RHIC and LHC is 0.9 and 4 respectively.

In hadronic collisions, one usually defines the "K-factor" as a measure of the size of higher-order corrections. Here we define the *effective* K-factor for *nuclear* collisions as a ratio of the particular distribution to the leading-order distribution without any nuclear effects. We find that in hadronic collisions K-factor is about 2 (at RHIC) and 3 (at LHC). The nuclear shadowing effect suppresses production of charm quarks by about 30% (at RHIC) and 53% (at LHC), giving the effective, *nuclear* K-factor of about 1.4 at both energies, in the central rapidity region ($|y| \leq 3$).

In Ref. 14 we present our results for the transverse momentum distribution of the charm quark produced in Au-Au collisions at RHIC and LHC (solid line). We find that both higher-order correlations and the nuclear effects change the shape of p_T distribution. At RHIC, the nuclear shadowing effect is much stronger at low p_T (about 40% effect), while at $p_T = 6GeV$ it reduces the cross section by only few percent. The next-to-leading order corrections give a factor of 1.7 increase at low p_T and about factor of 3 at $p_T = 6GeV$. These two effects together result in effective K-factor increasing from 1 at $p_T = 1GeV$ to about 3 at $p_T = 6GeV$. At large p_T, where nuclear shadowing effects are negligible, we expect K-factor to approach its hadronic value. At LHC, the nuclear shadowing effect is about 60% at low p_T and 40% at $p_T = 6GeV$. The nuclear K-factor increases from 0.4 at $p_T = 1GeV$ to about 7 at $p_T = 6GeV$. We find similar behavior of the K-factor for x_F distribution, namely its strong dependence on x_F.

1.1.3 Total and Single Differential Cross Sections for Charm Production

By integrating differential distributions over the phase space we obtain the total number of charm quark pairs produced. For central Au-Au collisions at RHIC (LHC) we get about 4 (40) charm quark pairs produced per event [5].

To summarize, we have calculated the complete next-to-leading order differential and total inclusive cross sections for the charm production in hadronic and nuclear collisions at RHIC and LHC energies [5,14]. We have shown that both higher-order contributions and the nuclear shadowing effect are large and can not be neglected. In the central region of the rapidity distribution, the higher-order contributions increase the cross section by a factor of 2 (3), while the nuclear shadowing effect result in additional decrease of about 1.4 (2.1). These two effects together result in *nuclear* K-factor of about 1.4 at both energies in the central rapidity region. On the other hand, in case of the p_T distribution, the K-factor changes from 1 at $p_T = 1GeV$ to 3 at $p_T = 6GeV$ at RHIC and from 0.4 at $p_T = 1GeV$ to 7 at $p_T = 6GeV$ at LHC. We have found that at RHIC and LHC energies the dominant subprocess for charm production is gg fusion ($\geq 95\%$). Therefore, future measurements of charm production in p-p and A-A collisions at RHIC and LHC energies, in addition to being a test of perturbative QCD, could provide valuable information about presently unknown A-dependence and x-dependence of the gluon density in a nucleus.

Finally, we make a remark on the possibility of detecting a signal for the formation of quark-gluon plasma via enchanced charm production in heavy-ion collisions at RHIC. We have found that 0.9 open charm quark pairs per unit rapidity (in the central region) will be produced in central Au-Au collisions via hard parton-parton scatterings. This result seems to indicate that the initial temperature of the quark-gluon plasma formed in Au-Au collisions at RHIC would need to be unrealistically high [4] to overcome the QCD background. However, further theoretical and experimental work is needed in better understanding of the nuclear shadowing effects on the gluon density, before definite conclusion can be made. We discuss this problem in details in Section 2.2.

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1.2 Effective Field Theory of Multiparticle Production in Nuclear Collisions

(I. Sarcevic, T. Elze and H. Eggers)

1.2.1 Intermittency in High Energy Collisions

Recent experimental measurements of fluctuations in transverse energy in ultrarelativistic heavy-ion collisions have provided evidence that nuclear constituents scatter and produce particle coherently [1]. The observed fluctuations in the deposited energy and multiplicity are remarkably large and can not be described by an independent-collision model [2]. The same conclusion has been reached when multiparticle density fluctuations in different phase space regions have been studied via factorial moments, defined as

$$F_i(\delta y) = \frac{1}{M}\sum_{m=1}^{M}\frac{\langle n_m(n_m-1)\ldots(n_m-i+1)\rangle}{\langle n_m\rangle^i}, \qquad (1)$$

where M is the number of rapidity bins ($\delta y = Y/M$) and n_m is the number of particles in the m^{th} bin [3]. In the past few years, these moments have been measured for different targets and projectiles at energy of 200GeV/nucleon [4]. They were found to increase with decreasing bin size indicating nonstatistical fluctuations and being incompatible with the predictions of the standard classical hadronization models embedded in the existing Monte Carlo models [4]. In addition, the observed effect, sometimes referred to as the "intermittency effect", can not be accounted for by the superposition of independent nucleon-nucleon collisions, even when rescattering and geometrical effects are included [3].

The possibility of creating the new form of matter, the quark-gluon plasma, in high-energy heavy-ion collisions have inspired intensive theoretical work on identifying the unambiguous QGP signal. The observation of the unusually large multiparticle density fluctuations have created a new excitement in the field, especially as a possibility of pointing towards the onset of the phase transition from quark-gluon plasma to hadronic matter. Phase transitions in QCD at high temperatures are of general interest – they are directly relevant to cosmology, since such a phase transition occurred throughout the universe during the early moments of the big bang and a first order phase transition could have altered primordial nuclear abundances. Unfortunately, up to now there are no conclusive predictions for detecting the quark matter in heavy-ion collisions and there is no theory to describe the observed "intermittency" phenomenon [5].

1.2.2 Multiparticle Correlations in Heavy-Ion Collisions

Multiparticle correlations in three "dimensions" are usually measured in the following way. A given total interval $\Omega_{\rm tot} = \Delta Y\, \Delta\phi\, \Delta P$ is subdivided into M^3 bins of side lengths ($\Delta Y/M$, $\Delta\phi/M$, $\Delta P/M$). With n_{klm} the number of particles in bin (k,l,m) and $n^{[q]} = n!/(n-q)!$ the "vertical" factorial moment is

$$F_q^v(M) \equiv \frac{1}{M^3} \sum_{k,l,m=1}^{M} \frac{\langle n_{klm}^{[q]} \rangle}{\langle n_{klm} \rangle^q} = \frac{1}{M^3} \sum_{k,l,m=1}^{M} \frac{\int_{\Omega_{klm}} \prod_i d^3x_i\, \rho_q(x_1 \ldots x_q)}{\left[\int_{\Omega_{klm}} d^3x\, \rho_1(x)\right]^q}. \tag{2}$$

The second equality illustrates how the factorial moment can be written in terms of integrals of the correlation function [6]. The alternative "horizontal" factorial moment is often preferred for three-dimensional analysis; this form, while being much more stable, has the drawback that it depends on the shape of the one-particle distribution function ρ_1.

To measure true particle correlations, known as cumulants, the trivial background must be subtracted. The first two cumulants are

$$C_2(x_1,x_2) = \rho_2(x_1,x_2) - \rho_1(x_1)\rho_1(x_2) \tag{3}$$

$$C_3(x_1,x_2,x_3) = \rho_3(x_1,x_2,x_3) - \sum_{perm} \rho_1(x_1)\rho_2(x_2,x_3) + 2\rho_1(x_1)\rho_1(x_2)\rho_1(x_3).$$

By integrating these relations over each bin, one can derive equations for integrated cumulants $K_q^v = \int \prod dx_k C_q / (\int dx \rho_1(x))^q$ in terms of the above vertical factorial moments [7],

$$K_2^v = F_2^v - 1, \qquad K_3^v = F_3^v - 3F_2^v + 2 \quad \text{etc.} \tag{4}$$

Whenever there are no true correlations, these cumulants become zero. It was found that in the case of heavy-ion collisions, there are only two-particle correlations: while K_2 is positive, the values of K_3, K_4 and K_5 have been found to be consistent with zero [8]. First found in terms of one-dimensional rapidity data, this has been confirmed by measurements by NA35 in two and three dimensions [9]. Corresponding findings for other nuclei and energies were published before [10].

1.2.3 Effective Field Theory of Multiparticle Correlations

We have seen that particles produced in high-energy heavy-ion collisions exhibit only two-particle correlations, indicating that perhaps higher-order correlations are washed out by rescattering of the initially correlated particles. Presently, there is no theory that describes this phenomena. Recently, we have proposed a three-dimensional statistical field theory of density fluctuations which has these features [11,12]. This model was formulated in analogy with the Ginzburg-Landau theory of superconductivity. The large number of particles produced in ultrarelativistic heavy-ion collisions justifies the use of a statistical theory of particle production. The formal analogy with the statistical mechanics of a one-dimensional "gas" was first pointed out by Feynman and Wilson and was later further developed by Scalapino and Sugar [13] and many others [5]. The idea is to build a statistical theory of the macroscopic observables by imagining that the microscopic degrees of freedom are integrated out and represented in terms of a few phenomenological parameters and by postulating that this theory will eventually be derived from a more fundamental theory, such as QCD.

While in the G-L theory of superconductivity the field (i.e. the order parameter) represents superconducting pairs, in the particle production problem, the relevant variable is the density fluctuation. The "field" $\Phi(\vec{x})$ is a random variable which depends on the rapidity of the particle and its transverse momentum $p_\perp$ and it is identified with the density fluctuation.

Even though particles produced in high-energy collisions need not be in thermal equilibrium, one can still introduce a functional of the field Φ, $F[\Phi]$, which plays a role analogous to the free energy in equilibrium statistical mechanics. In principle one should be able to derive this functional from the underlying dynamics.

We define a random field Φ as a function in a three-dimensional space spanned by $(y, \phi, p_\perp)$. Throughout, $p_\perp$ will be implicitly divided by a constant scale $\mathcal{P}$ so that it is dimensionless. Since we are not looking for a phase transition, we omit the quartic term and start with the functional [11]

$$F[\Phi] = \int_0^P dy \int_{-P/2}^{P/2} d^2 p_\perp \left[a^2 \left(\partial\Phi/\partial y\right)^2 + a^2 \left(\nabla_{\vec{p}_\perp}\Phi\right)^2 + \mu^2\Phi^2 \right] . \tag{5}$$

Taking the appropriate functional derivative, we find for the functional (5) the three-dimensional form of the two-point function

$$\langle \Phi(\vec{x}_1)\Phi(\vec{x}_2)\rangle = \frac{1}{8\pi a^2}\frac{e^{-R/\xi}}{R} , \tag{6}$$

where $\xi = a/\mu$ and $R \equiv [(y_1 - y_2)^2 + p_{\perp 1}^2 + p_{\perp 2}^2 - 2p_{\perp 1}p_{\perp 2}\cos(\phi_1 - \phi_2)]^{1/2}$. Further, we define $\Phi(\vec{x})$ as the fluctuation at the point $\vec{x}$ of the particle density for a given event, $\hat{\rho}_1(\vec{x})$, above/below the mean single particle distribution ρ_1 at that point:

$$\Phi(\vec{x}) \equiv \frac{\hat{\rho}_1(\vec{x})}{\rho_1(\vec{x})} - 1 . \tag{7}$$

Through these definitions, we find that $\langle \Phi(\vec{x}_1)\Phi(\vec{x}_2)\rangle = k_2(\vec{x}_1, \vec{x}_2)$ and that all higher order cumulants become exactly zero, $k_{q\geq 3} = 0$. By means of the specific form of the functional (5) and the definition of Φ as a fluctuation, we take account of the experimental facts in this regard. What is not experimentally certain and is to be tested is whether the second order correlations obey the Yukawa form (6).

1.2.4 Projections of Multiparticle Correlations to Lower Dimensions

The second reduced cumulant $k_2 \propto e^{-R/\xi}/R$ can be compared to data only after a suitable integration over its variables. For three dimensions, the vertical integrated cumulant is given by $K_2^v(\delta y, \delta\phi, \delta p) = F_2^v - 1 = M^{-3}\sum_{k,l,m=1}^{M} K_2^v(k,l,m)$, (always taking $\vec{x} \equiv (y, \phi, p_\perp)$), with

$$K_2^v(k,l,m) = \frac{\int_{\Omega_{klm}} d^3\vec{x}_1\, d^3\vec{x}_2\; C_2(\vec{x}_1,\vec{x}_2)}{\left[\int_{\Omega_{klm}} d^3\vec{x}\;\rho_1(\vec{x})\right]^2} = \int_{\Omega_{klm}} d^3\vec{x}_1\, d^3\vec{x}_2 \frac{k_2(\vec{x}_1,\vec{x}_2)\rho_1(\vec{x}_1)\rho_1(\vec{x}_2)}{\left[\int_{\Omega_{klm}} d^3\vec{x}\,\rho_1(\vec{x})\right]^2} , \tag{8}$$

i.e. the integration of k_2 involves a correction due to the shape of the one-particle three-dimensional distribution function $\rho_1(\vec{x})$. Eq. (8) as it stands is exact; horizontal versions have also been derived [12]. A first test of our model would therefore be to see if Eq. (4) or its horizontal equivalent obeys the data in $(y, \phi, p_\perp)$.

The theoretical $k_2(\vec{x}_1, \vec{x}_2)$ is further tested by comparing to factorial cumulant data of lower dimensions. For example, in (y, ϕ), the cumulant is $K_2^v(\delta y, \delta\phi) = M^{-2}\sum_{lm} K_2^v(l,m)$

with $p_\perp$ integrated over the whole window ΔP,

$$K_2^v(l,m) = \int_{\Omega_m} dy_1 dy_2 \int_{\Omega_l} d\phi_1 d\phi_2 \int_{\Delta P} dp_1 dp_2 \, \frac{k_2(\vec{x}_1,\vec{x}_2)\,\rho_1(\vec{x}_1)\rho_1(\vec{x}_2)}{\left[\int_{\Omega_m} dy \int_{\Omega_l} d\phi \int_{\Delta P} dp \; \rho_1(\vec{x})\right]^2}. \qquad (9)$$

Cumulants of other variable combinations and lower dimensions are obtained analogously. With these relations it is thus possible, given any three-dimensional theoretical function k_2 (or r_2), to compute factorial cumulants and moments for any combination of its variables [12]. Doing this for different variables serves as a strong test of the theoretical function as the moments probe its different regions.

In Figures 2-3 (see Ref. 12) we present our results for the vertical and horizontal factorial moments. In our calculation of the projections, we have made the following approximations: We factorize the one-particle distribution into its separate variables: $\rho_1(\vec{x}) = \langle N\rangle_\Omega\, g(y)\, h(\phi)\, f(p_\perp)$, where the three distributions g, h and f are separately normalized over their respective total intervals ΔY, $\Delta\Phi$ and ΔP. The azimuthal distribution is taken as flat, $h(\phi) = 1/\Delta\Phi$. We use the full experimental parametrization for $f(p_\perp)$ provided by NA35. The choice of two parameters a and ξ in Figures 2-3 [12] are given only as an illustration. They should be determined by comparison with three-dimensional data. Once they are fixed, the two-dimensional and one-dimensional projections are genuine predictions of the theory.

In summary, we have constructed a three-dimensional effective field theory of multiparticle correlations, which gives no higer-order correlations, in agreement with the recent heavy-ion data. In our theory, two-particle correlations have Yukawa form and the corresponding integrated cumulants have singular behavior for small regions of phase space. This prediction seems to be in qualitative agreement with the recent NA35 data [14]. In addition, we have shown that once the parameters a and ξ are determined from comparison with three-dimensional data, our theory gives genuine predictions for the two-dimensional and one-dimensional cumulants.

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1.3 High Energy Photonuclear Interactions

(R. Gandhi, L. Durand, K. Honjo, H. Pi and I. Sarcevic)

1.3.1 Hadronic Structure of the Photon

In the past few years, the "hadronic" interactions of the photon have become a focus of many theoretical and experimental studies. In particular, recent photoproduction measurements at HERA energies [1] have provided important confirmation of the hadronlike character of the photon, the fact that photon can produce $q\bar{q}$ pair, and then through subsequent QCD evolution fill up the confinement volume with quarks and gluons with a density akin to that of a pion or a nucleon [2,3]. This virtual quark-gluon structure of the photon gives significant contribution to the inclusive jet cross section in high-energy γp collisions, which has recently been measured by ZEUS and H1 Collaborations [4]. One could also observe the hadronic behavior of the photon by measuring the total cross section. Even low-energy measurement of the total photoproduction cross section up to $\sqrt{s} = 18GeV$ show some rise with energy [5] similar to the one observed in hadronic collisions [6]. This is even more apparent from the recent measurements of $\sigma^{\gamma p}$ at HERA energies [1]. In hadronic collisions, the rapid growth of the total cross section was associated with a dominance of hard scattering partonic processes over the nonperturbative (soft) ones [7], supported by the detection of the semi-hard QCD jets (so-called "minijets") at CERN Collider energies [8]. Similarly recent observation of the hard scatterings in photoproduction at HERA energies corroborates this hypothesis [5]. The hadronic interaction of the photon in ultra-high energy photon-nucleus collisions has also been studied [9] in the context of anomalously high muon content observed in PeV cosmic-ray air showers (the so-called "muon puzzle") [10]. The reported muon content was about order of magnitude higher than what one would expect from the electromagnetic showers initiated by the photon. However, if there is some threshold energy at which the photon becomes hadronlike, the photonuclear cross section would increase with energy in a similar way to the hadronic cross section giving larger probability for producing muons than the "conventional" value. We find that the ultra-high energy photonuclear cross section obtained in the QCD model [3] with realistic photon structure function and including the appropriate hadronic probability of the photon is too small to account for anomalies, but large enough to be interesting for study at HERA and Fermilab energies.

1.3.2 Jets, Jet Multiplicities and the Total Photoproduction Cross Sections

The low-energy measurements of the total photoproduction cross cross section in the energy range $10GeV \leq \sqrt{s} \leq 18GeV$ [5] and, especially, recent HERA data [1] point towards the hadronic behavior of the photon. Few years ago, we have made predictions for the total and jet photoproduction cross sections in a simple QCD minijet-type model based on analogy with hadronic collisions [2]. We have assumed that the total photoproduction cross section can be represented as a sum of the soft (nonperturbative) and hard (jet) part (i.e. $\sigma_T = \sigma_{soft} + \sigma_{JET}$), where the soft part is energy independent, determined from the low energy data ($\sqrt{s} \leq 10GeV$). The jet (hard) part has contributions from two subprocesses: the "standard" (direct) QCD process ($\gamma q \rightarrow qg$ and $\gamma g \rightarrow q\bar{q}$) and the "anomalous" process (for example, $\gamma \rightarrow q\bar{q}$, followed by quark bremsstrahlung, $q \rightarrow qg$ and $gg \rightarrow gg$). The later process is the same as the jet production process in p-p collisions up to a photon structure function. We note that the photon structure function is proportional to α_{em}/α_s, where α_{em} is the electromagnetic coupling. The effective order of the above processes is therefore $\alpha_{em}\alpha_s$, since the jet cross sections are of order α_s^2. Thus, they are of the same order as direct two-jet processes, in which the photon-parton vertex is electromagnetic and does not involve the photon's hadronic content. The existing parametrizations of the photon structure function, Duke and Owens (DO), Drees and Grassie (DG) and Abramowicz et al. (LAC1) [11], all describe low energy photoproduction data very well. However, they differ dramatically at very high energies (i.e. low x region), the region of the HERA experiment, for example. Therefore recent photoproduction measurements at HERA energies can provide valuable information about the photon structure function at small x ($x \sim 4p_t^2/s$), and in particular its gluon content.

The QCD jet cross section for photon-proton interactions is given by [3]

$$\sigma^{\gamma p}_{\rm QCD} = \sum_{ij} \frac{1}{1+\delta_{ij}} \int dx_\gamma dx_p \int_{p^2_{\perp,\rm min}} dp_\perp^2 [f_i^{(\gamma)}(x_\gamma, \hat{Q}^2) f_j^{(p)}(x_p, \hat{Q}^2) + i \leftrightarrow j] \frac{d\hat{\sigma}_{ij}}{dp_\perp^2}, \qquad (1)$$

where $\hat{\sigma_{ij}}$ are parton cross sections and $f_i^{(\gamma)}(x_\gamma, \hat{Q}^2)$ ($f_j^{(p)}(x_p, \hat{Q}^2)$) is the photon (proton) structure function. The expressions for all the subprocesses that contribute to $\hat{\sigma_{ij}}$ can be found in Ref. [12]. We take the choice of scale $\hat{Q}^2 = p_T^2$, which is shown to give very good description of the hadronic jet data [12]. For parton structure function we use EHLQ parametrization [13]. The results do not show appreciately sensitivity to the choice of the proton structure function.

From the constant low energy data [5], we determine the soft part of the cross section to be $\sigma_{soft} = 0.114mb$. The observed 3% increase of the cross section in the energy range between $10GeV$ and $18GeV$ can be described by adding the hard (jet) contribution with jet transverse momentum cutoff $1.4GeV \leq p_{\perp,\rm min} \leq 2GeV$ to the soft part [2]. The actual value of $p_{\perp,\rm min}$, however, below which nonperturbative processes make important contributions, is impossible to pin down theoretically using perturbative techniques. As the energy increases direct and soft part become negligible in comparison with the anomalous part, because the later has a steep increase with energy. We find that in the Fermilab E683 energy range ($\sqrt{s} \leq 28 GeV$), the results for the cross sections are not sensitive to the choice of the photon structure function. Therefore, in addition to providing important confirmation of the hadronlike nature of photon-proton interactions, one could use forthcoming E683 experiment to pin down the theoretical parameter $p_{\perp,\min}$ to a few percent.

In Ref. 2 (Fig. 1), we present our results for $\sigma^{\gamma p}$ at HERA energies. We find results to be very sensitive to the choice of the photon structure function due to their different x behavior at very high energies. For example, DO gluon structure function behaves as $f_g^\gamma = x^{-1.97}$, while DG has less singular behavior, $f_g^\gamma \sim x^{-1.4}$ at the scale $Q^2 = p_T^2 = 4GeV$. The cross sections obtained using DG photon structure function are more realistic, since the extrapolation of DO parametrization to small x region give unphysically singular behavior. For this reason, all the cross sections obtained using DO function are only given as an illustration how $\sigma^{\gamma p}$ depends strongly on the choice of the photon structure function. In Ref. 3 we also present the results for the cross section when only soft and direct part are included, indicating its very weak energy dependence. The rise of the total cross section is thus mostly driven by the "anomalous" (hadronic) part of the cross section. We note that HERA measurement has some resolving power to distinguish between different sets of photon structure function and therefore determine presently unknown low x behavior of its gluon part. For example, the cross sections obtained using DO structure functions are already excluded by HERA data, while theoretical result obtained using DG structure function and $p_{\perp,\min} = 2GeV$ is consistent with the data. However, one should keep in mind that all the theoretical predictions presented in Ref. 2 do not take into account the possibility of multiple scatterings which, as we show later, result in reducing the cross sections by 10% for $p_{\perp,\min} = 2GeV$ and by about 25% for $p_{\perp,\min} = 1.4GeV$ [3].

We have also calculated the total jet cross sections at Fermilab and HERA energies for jet p_T triggers of $3, 4, 5GeV$ and $5, 10, 15GeV$ respectively [2]. The energy dependence of the total jet cross section is much steeper than of the direct part of the jet cross section only. This can be seen in Figs. 4 and 5 in Ref. [2]. Again, jet measurements at HERA energies can distinguish between different photon structure functions, but in this case one is probing the photon structure function at slightly higher values of x than in the case of the total cross section.

Recently, we have improved our calculation of the total photoproduction cross section by incorporating the possibility of multiple parton collisions in a single photon-nucleon collision by using an eikonal treatment of the high-energy scattering process [3]. In our general formulation of the theory the γp cross section is expressed as a sum over properly eikonalized cross sections for the interaction of the virtual hadronic components of the photon with the proton, with each cross section weighted by the probability with which that component appears in the photon. We have developed a detailed model which includes contributions from light vector mesons and from excited virtual states described in a quark-gluon basis. The parton distribution functions which appear can be related approximately to those in the pion, while a weighted sum gives the distribution functions for the photon. The "soft" contributions to the eikonal functions are parametrized in the form expected from Regge theory to assure that the calculated high-energy cross sections connect smoothly with the cross sections measured at lower energies. The parton distribution functions which appear in the "hard scattering" contributions to the eikonal functions can be related approximately to those in the pion. This approach allows us to use the (not-well-determined) pion distributions to predict the high-energy behavior of the γp cross section, and also gives an approximate prediction for the parton distributions in the photon in terms of these in the pion. We have also developed an alternative approach based directly on the photon distribution functions, in accord with the analysis done in Ref. [2].

In our QCD-based eikonal model the total inelastic γp cross section is given by [3]

$$\sigma_{\rm inel}^{\gamma p} = \sigma_{\rm dir} + 2\lambda \mathcal{P}_\rho \int d^2b \left(1 - e^{-{\rm Re}\,\chi_{\rho p}}\right) + 2\sum_q e_q^2 \frac{\alpha_{\rm em}}{\pi} \int_{Q_0^2} \frac{dp_{\perp 0}^2}{p_{\perp 0}^2} \int d^2b \left(1 - e^{-{\rm Re}\,\chi_{q\bar{q}p}}\right), \quad (2)$$

where

$$2{\rm Re}\,\chi_{\rho p}(b,s) = A_{\rho p}(b)\left[\sigma_{\rm soft}(s) + \sigma_{\rm QCD}^{\rho p}(s)\right], \quad (3)$$

and

$$2{\rm Re}\,\chi_{q\bar{q}p}(b,s,p_{\perp 0}) = A_{q\bar{q}p}(b,p_{\perp 0})\left[\sigma_{\rm soft}^{q\bar{q}p}(s,p_{\perp 0}) + \sigma_{\rm QCD}^{q\bar{q}p}(s,p_{\perp 0})\right]. \quad (4)$$

$A_{\rho p}$ and $A_{q\bar{q}p}$ are the overlapping spatial densities of the corresponding systems and $\sigma_{\rm QCD}^{q\bar{q}p}$ is given by the analog of Eq. (1) with f_i^ρ replaced by $f_i^{q\bar{q}} \approx (Q_0/p_{\perp 0})^2 f_i^\rho$. In Eq. (2), the parameter λ is the "weight" for the low-mass vector meson contributions. For example, $\lambda = 4/3$ for equal ρp, ωp and ϕp cross sections and $\lambda = 10/9$ for complete suppression of the ϕ contribution.

The incident hadronic systems in a ρp collision can interact inelastically through soft as well as hard processes. The soft scattering is dominant at low energies. Motivated by Regge theory, we parametrize $\sigma_{\rm soft}$ as energy-dependent, i.e.

$$\sigma^{\rho p}_{\rm soft}(s) \approx \sigma_0 + \sigma_1 (s - m_p^2)^{-1/2} + \sigma_2 (s - m_p^2)^{-1}. \tag{5}$$

We take $\sigma^{\rho p}_{\rm QCD}$ to be equivalent to $\sigma^{\gamma p}_{\rm QCD}$ given by Eq. (1). We find that QCD contributions to the hadronic interactions of the photon still lead to a rapid rise in $\sigma^{\gamma p}$ at HERA energies as predicted in earlier calculations [2] and observed in recent experiments [1]. The magnitude of the rise provides a quantitative test of the whole picture. In particular, our results presented in Fig. 2 in Ref. 3, show clearly that measurements of the total inelastic γp cross section at HERA can impose strong constraints on the parton distributions in the photon and, when combined with low-energy measurements, it can pin down the value of the theoretical cutoff $p_{\perp,\rm min}$, used to determine the onset of hard-scattering processes. We find that the cross sections obtained using the value of the cutoff $p_{\perp,\rm min} = 2GeV$ seem to be a bit too low for the observed increase of the cross section in the energy range between $10GeV$ and $18GeV$, but is in excellent agreement with ZEUS and H1 data, while our results with $p_{\perp,\rm min} = 1.41GeV$ are in better agreement with the data at all energies when pion structure function is used and slightly too large at HERA energies when DG structure function is used for the photon. We find that the eikonalization effect reduces the cross sections at HERA energies by about 10% for the case of $p_{\perp,\rm min} = 2GeV$ and by almost 25% when $p_{\perp,\rm min} = 1.4GeV$ is used.

1.3.3 Ultrahigh Energy Photonuclear Cross Section and the “Muon Puzzle”

The unusually large photonuclear cross sections at very high energies play an important role in understanding recently observed anomalous muon content in cosmic ray air-showers associated with astrophysical “point” sources (such as Cygnus X-3, Hercules X-1 and Crab Nebula) [10]. The number of muons observed is comparable with what one would expect in a hadronic shower, but the fact that primary particle has to be long-lived and neutral, makes photon the only candidate in the Standard Model. Conventionally, one would expect that photon produces electromagnetic cascade and therefore muon poor. However, if the photonuclear cross section at very high energies becomes comparable with pair production and bremsstrahlung cross section ($\sigma_{\gamma\to e^+e^-} \sim 500mb$) the muon content in a photon initiated shower will be affected. The hadronic character of the photon enchances the photonuclear cross section at very high energies [2]. We have calculated the total inelastic photon-air cross sections in a QCD-based diffractive model, which takes into account unitarity constraints necessary at ultra-high energies [3].

Our model for the hadronic interactions of the photon can be extended to photon-nucleus interactions using Glauber's multiple scattering theory with the result [3]

$$\sigma_{\rm inel}^{\gamma-{\rm air}} = A\sigma_{\rm dir} + 2\lambda\mathcal{P}_\rho \int d^2b\, \langle\Psi|1 - \exp\left(-\sum_{j=1}^{A} {\rm Re}\,\chi_j^{\rho p}\right)|\Psi\rangle +$$

$$+2\sum_q e_q^2\frac{\alpha_{em}}{\pi}\int_{Q_0^2}\frac{dp_{\perp 0}^2}{p_{\perp 0}^2}\int d^2b\,\langle\Psi|\,1 - \exp\left(-\sum_{j=1}^{A}{\rm Re}\chi_j^{q\bar{q}p}\right)|\Psi\rangle. \qquad (6)$$

The expectation values are to be taken in the nuclear ground state. ${\rm Re}\chi_j^m = {\rm Re}\chi^m(|\mathbf{b}-\mathbf{r}_{j\perp}|)$ is the eikonal function for the scattering of the hadronic component $|m\rangle$ of the photon on the $j^{\rm th}$ nucleon, where $\mathbf{r}_{j\perp}$ is the instantaneous transverse distance of that nucleon from the nuclear center of mass, and $\mathbf{b}$ is the impact parameter of the photon relative to the nucleus. We have evaluated the integrals using shell-model wave functions for the oxygen and nitrogen nuclei.

Our results for the photon-proton and photon-air cross sections at energies $10GeV \leq \sqrt{s} \leq 10^4 GeV$ are presented in Fig. 6 in Ref. [3]. We find that $\sigma^{\gamma p}$ at $\sqrt{s} \geq 10^3 GeV$ is about twice the conventional value, large enough to be interesting, but much too small to account for the reported muon anomalies in photon-initiated showers.

In brief, we have shown how measurement of the total and jet photoproduction cross section at HERA energies provides valuable information about the hadronic character of the photon and, in addition, it can impose strong constraint on the value of the theoretical jet momentum cutoff and, more importantly, determine the low x behavior of the photon structure function and its gluon content. With the theoretical uncertainties being reduced, we have extrapolated our predictions for the photonuclear cross section to ultra-high energies relevant for cosmic ray experiments. The new results on γ-air interactions make it quite clear that the hadronic interactions of the photon cannot explain the reported muon anomalies in cosmic ray air-showers, if the anomalies in fact exist. Furthermore, future cosmic ray experiments might be able to put the "muon puzzle" observations on firmer grounds and to provide valuable input to particle physics at ultra-high energies, currently far beyond the range of accelerator experiments.

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1.4 Rapidity Gaps in High-Energy Collisions

(R. Fletcher and T. Stelzer)

Rapidity gaps are regions of rapidity in which no hadrons are produced in a high energy hadronic interaction. The production of these gaps has received much attention as a possible signature for new physics, especially for the production of Higgs Bosons [1,2]. Recently we have proposed a rapidity gap analysis for studying multiparticle production in high energy collisions [3]. In one considers events characterized by two high p_t jets widely separated in rapidity, the main contribution in the production of these jets is described by tree level QCD. The partons, quarks or gluons from the beam hadrons, exchange a gluon at high momentum transfer, and produce jets. A second contribution comes from the exchange of a color singlet between the incoming partons. This can be a photon or weak boson, or a combination of gluons which form a color singlet, the hard pomeron. In each case the outgoing partons radiate, and fragment, forming high p_t jets in the detector. In addition to the hard jets in a hadronic interaction, there is an 'underlying event' of soft particles which typically has the structure of a minimum bias event, and there are particles generated by the fragmentation of the soft partons radiated from the hard scattering.

The production of the underlying event is very similar to particle production in minimum bias events [4]. For instance, when a two jet system is formed, the beam particles 'lose' a parton to the 2-jet system, so the beam fragments carry color. In order to form a color singlet the forward and backward beam fragments must be joined by a string. Again, this fills the central rapidity region with particles. Minijets appear again in the underlying event. Secondary interactions between partons can produce minijets underlying the main, high p_t reaction.

Jet events caused by color singlet exchange have a different structure. In these events the incoming quarks exchange a color singlet boson, or pomeron, so the recoiling quark which forms a forward jet, and the surviving beam remnant form a color singlet. Since these partons are both in the forward (backward) region, they won't create particles in the central region.

Equally important is the fact that parton, generally gluon, radiation in the region between the tagging jets is suppressed [5]. The difference in the radiation in QED and QCD exchange events can be understood from a simple classical analogy. From electrodynamics, we know that accelerated charges radiate, and tend to radiate collinearly. In QCD when a color charge is accelerated, it will result in radiation, primarily tangentially to the path of the accelerated charge. For color singlet exchange, the color is accelerated through a small angle, and radiation is generated only in a small region of phase space. For color octet exchange, even if the incoming fermions scatter through a smal angle, the color charge accelerates through a large angle, because the charge is exchanged between the scattering partons. This produces radiation over a large region of rapidity.

We can see this effect in perturbation theory by calculating the 2 to 3 matrix elements, and plotting the ratio of the QED to QCD contributions for fixed quark jets, while varying the rapidity of the radiated gluon. If we plot the ratio of the QCD to QED exchange matrix elements, varying the rapidity of the radiated gluon, we find that there is, as expected, a strong supression of radiation into the gap region [5,6].

To summarize, radiation is suppressed in the central region when a color singlet is exchanged between partons. This effect can be understood classically in terms of 2 to 3 Feynman amplitudes. Rapidity gaps recently observed by CDF and D0 at the Tevatron have confirmed our predictions [7].

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1.5 High Energy Nuclear Interactions in Cosmic Ray Cascades

(R. S. Fletcher, T. Gaisser, F. Halzen and T. Stanev)

Dr. Fletcher, our former postdoctoral fellow, has developed a detailed Monte Carlo simulation of particle production in hadron-hadron, and hadron-nucleus interactions [1]. Nuclear interactions are included by using Glauber's wounded nucleon expansion. A principle application of the program is to simulate high energy multiparticle production in very high energy cosmic ray showers. Here, the target is nitrogen, or oxygen. The program has been interfaced to the HEMAS and MOCCA air shower simulations. Several cosmic ray collaborations are evaluating the program for possible use in their data analysis.

The present SIBYLL model treats hadron-nucleus collisions. Preliminary version of the program which simulates heavy ion collisions has been recently developed [2]. This program is the first simulation that is able to treat nuclear effects in high energy collisions in an air shower simulation. SIBYLL and most other models of high energy multiparticle production are based on minijet production and string fragmentation. These models are poor at reproducing the multiparticle correlations seen in experimental data and need to be improved.

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1.6 Bose-Einstein Correlations in High Energy Collisions

(P. Carruthers, P. Lipa, M. Greiner and H. Eggers)

1.6.1 Influence of Bose-Einstein Correlations on Intermittency

In recent years the intermittency phenomenon has been well established in multiparticle production and was confirmed by many experimental groups. While in e^+e^- physics the rise of factorial moments with decreasing bin size is mostly attributed to the cascade like branching of energetic partons in jets, the situation in soft hadron-hadron and nucleus-nucleus collisions is still far from being understood.

One contribution to intermittency in soft physics are the so-called Bose-Einstein (BE) correlations. Therefore we investigated the influence of these well known but not well understood short range correlations on the intermittency phenomenon in $\bar{p}p$ reactions at $\sqrt{s} = 630\,\mathrm{GeV}$. This work was done in close collaboration with B. Buschbeck and N. Neumeister from the Institute for High Energy Physics within the framework of the UA1 Minimum Bias Collaboration.

In our work [1] we provided the tools that makes a unified Bose-Einstein and intermittency analysis possible; the commonly used factorial moments F_q are not well suited for such a unified study since they are based on absolute coordinates whereas BE correlation measurements require the use of relative coordinates (three or four momentum differences between like sign particles). A generalization of the commonly factorial moments F_q using relative coordinates — the correlation integrals C_q — allows for a formulation of the intermittency phenomenon and BE correlations in terms of the same observables and thus a common study of both effects.

We used two different kinematical variables for this study: a) the 1-dimensional pseudorapidity difference $\delta\eta$ between two particles and b) the 4-momentum difference $Q^2 = -(p_1^\mu - p_2^\mu)^2$, which is commonly used in the BE correlation analyses. For higher orders $q > 2$ a generalization $Q^2 = q_{12}^2 + q_{23}^2 + \cdots + q_{(q-1)q}^2$ was used which is simply related to the invariant mass M of the q-tuple by $Q^2 = M^2 - (qm_\pi)^2$ where m_π is the pion mass.

The conclusions of this work are:

The influence of the BE effect on intermittency depends strongly on the variable used. Whereas it is weak in the case of $\delta\eta$, it is the dominant contribution to the rise of C_q with decreasing $Q^2 \to 0$. A comparison of slopes φ_q obtained by fits to the power law $C_q(Q^2) \propto (Q^2)^{-\varphi_q}$ for a like–sign (ls) sample and a sample consisting of all charged particles (all) reveals that there is no significant difference in slopes ($\varphi_q^{all} \approx \varphi_q^{ls}$) for the $\delta\eta$–analysis; however, in Q^2 variables the slopes differ by a factor of two, i.e. $\varphi_q^{all} \approx \varphi_q^{ls}/2$. This is fullfilled for all experimentally accessible orders $q \leq 5$.

Different dynamical mechanisms are dominant in the $\delta\eta$ and Q^2 analyses, which was shown by a direct comparison of the two particle correlation function of like–sign pairs to the one of unlike–sign pairs: when analyzed in $\delta\eta$, the two particle correlation function is dominated by the contribution of unlike–sign pairs whereas in Q^2 the dominance of like–sign pair correlations shows up very clearly for small $Q^2 < 0.03\,\mathrm{GeV}/c$.

The intermittency analysis in Q^2 — carried out for the first time in this work — shows an almost perfect power law dependence of C_q on Q^2 over the whole region of analysis. It is due to the interplay of different mechanisms, however, for small Q^2 only the BE effect remains (at least in the 2–body case). Therefore the question, if intermittency in soft hadron-hadron physics can be reduced to known effects will be answered probably with "yes"; but we stress the fact that this investigation also brings a completely new aspect into the understanding of the BE effect: a very steep power law dependence on Q^2 emerges and we do not see any cut–off mass yet. Such a power law — if confirmed by further studies — is a new and unexpected feature of BE correlations and contradicts previous parametrizations by Gaussians or exponentials. According to the discussion in [2], this could suggest a completely irregular shape of the source which extends sometimes to very large distances ($\approx 6fm$). It could indicate that the interaction region itself is a fractal.

Further experimental and theoretical studies are required to confirm the power law dependence of the correlation integrals on Q^2 and to understand its physical origin. The concept of a fractal source is new and very attractive, but other effects, such as the contribution from long living resonance decays to the BE effect, have to be understood and separated first.

1.6.2 Measurement of True Higher Order Bose-Einstein Correlations

Lately, theoretical and experimental work in pion interferometry has shifted towards the study of higher order Bose–Einstein correlations, i.e. correlations between more than two particles. The correlation integral technique introduced by us [1] made it possible to measure moments C_q in dependence of Q^2 up to order $q = 5$. But it is well known that higher order moments contain a large combinatorial background from lower order ones. Thus, it is an open question whether true higher order correlations exist in pion interferometry or whether they are negligible and of no further theoretical and experimental interest.

While this is due partly to limited statistics of experimental data, it is also a consequence of the fact that previously no clear recipe how to access true higher order correlations (i.e. the higher order cumulants) had been available.

In our recent work [3] we have developed a consistent framework to measure cumulants as a function of Q^2 and/or the invariant mass M of a system of particles. By putting the previously heuristic technique of event mixing on a sound footing, we have justified its use. Extending this technique to combinations of correlated and uncorrelated terms in the expansion of cumulant functions (e.g terms of the form $\rho_1(x_1)\rho_2(x_2,x_3)$), we have shown how the combinatoric background may be eliminated. We have little doubt that this background is large; it remains to be seen whether, after its elimination, any residual "true" correlations are left.

We have to stress that the simple expansions of cumulants in terms of normalized factorial moments (e.g. $K_3 = F_3 - 3F_2 + 2$) are not valid for correlation integrals in Q^2. The normalized cumulants $K_q(Q^2)$ must rather be calculated from event mixing structures, otherwise spurious correlations may be introduced. It is possible that such naive expansions are valid approximations of the exact forms, but this cannot be assumed a priori and must be tested in further studies. The lesson is that the cumulant expansions must be made at the primary level of the three–momenta correlation functions before any integration and/or projection is carried out.

Our work is restricted for the time being to the conceptual and combinatorial aspects of the experimental measurements of true higher order correlations of like–sign as well as all charged particles in Q^2. The intricate and murky question of separating correlations due to Bose statistics from other kinematical or dynamical correlations has to be postponed to forthcoming studies.

1.7 Refined Measurement of Higher Order Correlations and Intermittency

(P. Carruthers, P. Lipa, H. Eggers and B. Buschbeck)

We have developed a general formalism for measuring correlations of point distributions. The language used has been that of high energy physics, but we believe that the method may be of use in other fields also.

Previously, we introduced a form of the correlation integral, which we named the *Snake integral* and compared it to a symmetrized version, the "GHP" integral [1]. In our more recent work [4], we advocate a slightly different form, which for obvious reasons we name "Star integral".

Generally, every correlation integral of order q assigns a size to every possible q-tuple of particles in an event. The way this assignment is done is not unique and distinguishes the different versions of correlation integrals: the Snake form takes the size of the longest segment of a line (snake) connecting the q particles, the GHP form takes the longest of all possible $q(q-1)/2$ interparticle distances within the chosen q-tuple, and the Star prescription assigns the longest distance of $q-1$ particles from a chosen center particle. Another form useful for pion interferometry, discussed in the previous sections, takes the invariant mass of the q-tuple as measure of its size. Once such an assignment is made, all the different correlation integrals count the abundance of q-tuples in dependence on their size.

The Star integral shares all the advantages of the Snake and GHP forms but is more accessible to intuitive understanding and, what is most important for the practical implementation, is computationally orders of magnitudes more efficient. The previous forms required a CPU time of order N^q (for the Snake) or $N^q/q!$ (for the GHP), where N is the total multiplicity of an event and q the order of the moment. Higher order correlation measurements were thus restricted by prohibitive CPU time consumption. The algorithms derived from the Star integral require only CPU time of order N^2, independent of the order

q of the moment! This means that this correlation integral can now be used also for correlation analyses in *heavy ion collisions*, something hitherto impossible due to the large event multiplicities involved. The big improvements in statistics over the conventional vertical factorial moments will allow for much more accurate measurements.

Furthermore, our formalism leads naturally to a conceptual cleanup of a heuristic form used in the characterization of strange attractors [5] and of scaling in galaxy distributions [6] in the sense that *factorial* powers rather than ordinary ones have to be used. The difference between these two becomes significant when small particle numbers are involved, a situation which occurs inevitably when the integration domain becomes sufficiently small.

Another objective of our work is to provide a technique to measure integrals of cumulant correlation functions, which are the genuine higher order correlations and thus very useful observables. Previously, the only way to extract information on higher order cumulants had been via combinations of factorial moments; here we show how cumulant functions can be integrated directly, by clarifying and extending the technique of event mixing, thus making use of all the advantages provided by correlation integrals over bin-averaged factorial moments.

The present work has emphasized the *how to* rather than the *what* of measurement. At first sight, it may seem unnecessary to devote so much effort to the mere process of measurement. However, high energy hadronic collisions, and even more so nucleus-nucleus collisions, are presently in a state where very few exact calculations on correlations can be done and most theoretical work relies on assumptions which often are hard to support or believe. In this context, we believe that it is of cardinal importance that there should be a clear and unambiguous method for analyzing correlations, and that, if possible, a standard should be established by which different experiments can be compared.

Three issues have been dealt with only cursority or not at all: the question of eliminating the influence of the total multiplicity distribution, the measurement of correlations between different species of particles (e.g. distinguished by their charge), and the problem of properly defining fractal dimensions for nonstationary ensembles of events as an extension of commonly used definitions in the study of chaotic single-event time series and strange attractors. We plan to address these issues in future work.

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1.8 Wavelet Correlations in Multiparticle Cascades

(P. Carruthers, M. Greiner and P. Lipa)

In the past ten years a powerful new mathematical method called wavelet transformation has been developed [1]. In contrast to the Fourier transform the wavelet transform is more or less local in both time and frequency and possesses a multiresolutional character. This resulted in powerful applications of the wavelet transformation in signal analysis and data compression. These properties of the wavelet transformation are also very attractive for the analysis and description of complex multiparticle cascades.

In a first step we discussed the wavelet transformed correlations of some discrete, hierarchical cascade models (p- and α-model, QCD-motivated cascade model); these models serve as simplified approximations to more realistic QCD parton cascade models. It turns out that the wavelet transformation (quasi-) diagonalizes the covariance matrix, which makes it extremely simple to read of selfsimilarity (and deviations thereof) of the underlying cascade process. This result reflects the compression power of wavelets! The other wavelet property, i.e. locality in conjugate variables, enters into higher order wavelet correlations. Those can be interpreted as interscale cluster correlations between structures (jets) living on different scales. Since these structures have been produced during different stages of the cascade process, these higher order wavelet correlations provide extremely sensitive information on the reaction dynamics.

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The following publications resulted so far:

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1.9 Other Members of the Group (1993-1995)

In the period of 1993-1995, Drs. Fletcher, Lipa and Gandhi were postdoctoral fellows at the University of Arizona. Drs. Greiner and Elze were Visiting Associate Professors at Arizona. Except for Prof. Gandhi, who is currently holding faculty position in India, the others are postdoctoral fellows at various academic institutions. Peter Valerio and Igor Zubarev are currently graduate students working under the supervision of Prof. Sarcevic. Dr. Zheng Huang has joined our research group since August 1995 as a postdoctoral fellow.

Chapter 2

Research Proposal

2.1 Heavy Quark Production at RHIC and LHC Energies to $O(\alpha_s^3)$

(I. Sarcevic and P. Valerio)

Recently, there has been a considerable theoretical and experimental interest in studying heavy-quark production in hadronic and nuclear collisions. Theoretical results for the next-to-leading order, $O(\alpha_s^3)$, radiative corrections of the parton level cross section [1] have been used to improve calculation of the total heavy-quark cross sections in hadronic collisions at energies up to 2TeV [2]. For bottom production these corrections are found to be large, especially at threshold energies and at very high energies [2]. Studying charm and bottom production at the proposed heavy-ion colliders, BNL's Relativistic Heavy Ion Collider (RHIC) and CERN's Large Hadron Collider (LHC) is of special interest. Open charm production has been suggested as an elegant signal for detecting the formation of quark gluon plasma (QGP) in heavy-ion collisions [3]. If thermalization is reached in heavy-ion collisions at RHIC ($\sqrt{s} \approx 200\text{GeV}/A$) and LHC ($\sqrt{s} \approx 7\text{TeV}/A$) energies we expect a very dense matter, of the order of few GeV$/fm^3$, to be formed in the initial stage of the collision. It seems plausible that this density is sufficient for creating a new state of matter, the QGP [4]. The search for a clean, detectable signal for this new state of matter has been one of the most challenging theoretical problem for the last few years. Some of the proposed signals, thermal photons, dileptons and J/Ψ suppression, have been studied and found to be difficult to detect due to the large QCD background [5]. In order that the enhanced charm production can be used as the signal of QGP, we need to understand the QCD background, namely the production of charm quarks through the hard collisions of partons inside the nuclei. This type of charm production at RHIC and LHC energies is dominated by initial-state gluons. Therefore, in addition to the possibility of pointing towards the formation of quark-gluon plasma in high-energy heavy-ion collisions, combined measurements of charm production in p-p, p-A and A-A collisions could provide valuable information about the gluon density in a nucleus.

In the past year, we have completed the full next-to-leading order calculation of the open charm production in hadronic and heavy-ion collisions at RHIC and LHC energies [6,7]. The total cross section for charm production in hadronic collisions was obtained by convolving the parton densities in the hadron with the hard scattering cross section. We have included the leading order subprocesses, $O(\alpha_s^2)$, and all the next-to-leading order subprocesses, $O(\alpha_s^3)$, such as $gg \to gQ\bar{Q}$, etc. In these calculations the factorization scales in the structure functions was taken to be $2m_c$ and the renormalization scales considered were m_c and $2m_c$. We took the charm mass to be $m_c = 1.5GeV$. The factorization scale was not taken to be below $2m_c$ because the structure functions have been measured only for $Q^2 \geq 8.5GeV$ [8]. We have used the two-loop-evolved parton structure functions of Martin, Roberts and Stirling: MRS D0, MRS D–, and MRS A. These distributions have different behavior at small-x. MRS D– [9] and MRS A [10] exhibit "singular" behavior at small-x, i.e. $G(x) \sim x^{-1.5}$ and $G(x) \sim x^{-1.08}$ respectively, while MRS D0 is not singular at small-x. The most recent HERA data [8] seems to be steeper than the MRS D0 set and not as steep as the MRS D– distributions. If the saturation effect of partons is included, with the assumption that partons are localized in "hot spots" inside the proton, the behavior of the MRS D– parton distributions is in agreement with HERA data. The x-range probed in the charm production at RHIC energies in any region of the phase space is well within the range of the HERA data, allowing for the more reliable theoretical predictions than for LHC energies. At the LHC, charm production in the forward rapidity region is sensitive to x-region order of magnitude smaller than the current HERA data. In order to reduce theoretical uncertainty due to the choice of the renormalization scale we have compared our total cross section calculation to the current data for charm production in pp and pA collisions for beam energies ranging from $50GeV$ to $2TeV$ [11]. We have found better agreement with the data when we used the renormalization scale $Q^2 = m_c$. The uncertainty due to the choice of the structure function for this relatively low energy region was negligible.

To determine charm production in nuclear collisions we need to know the parton densities *in-medium*. If nucleons were independent, the parton densities in a nucleus would simply be A times the parton density in a nucleon. At high energies, the parton densities become

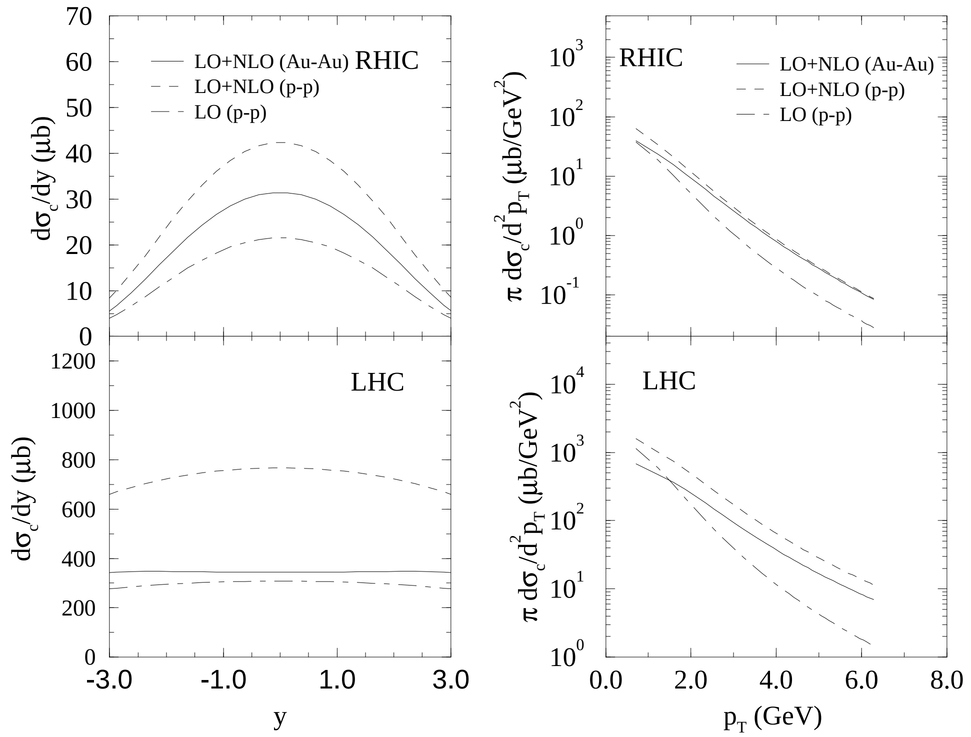

Figure 2.1: The rapidity and p_T distributions for charm quark production in p-p and Au-Au Collisions at RHIC and LHC energies. The Au-Au cross sections are scaled by $1/A^2$ so that a comparison to p-p collisions can be made. These distributions are calculated to $O(\alpha_s^3)$ with MRS A structure functions.

so large that the sea quarks and gluons from different nucleons overlap spatially and the nucleus cannot be treated as a collection of uncorrelated nucleons. When partons from different nucleons begin to overlap they interact and through annihilation effectively reduce the parton densities in a nucleus. From the EMC, NMC, and E665 experiments [12] we see that the parton densities per nucleon are effectively reduced at small-x for large A. This effect has been parameterized in terms of shadowing function, $R(x, A)$. This function, when multiplied by the nucleon parton densities reproduces the observed suppression at small-x.

Measurements of the nuclear shadowing effect [11] imply that the shadowing function has strong x-dependence. This indicates that the effect that nuclear shadowing has on charm production is not uniform throughout phase space. Also, the new processes that are part of the next-to-leading order calculation have different contributions in different regions of phase space. Thus, we have also calculated the single inclusive distributions, such as the rapidity and transverse momentum distributions, as well as the double differential distributions. We present our results for representative sets of distributions in Fig. 2.1. We note that shadowing effect has strong p_T dependence. Furthermore, the is stronger at LHC energies where lower x region is probed. To illustrate the importance of the $O(\alpha_s^3)$ corrections *combined* with nuclear shadowing effects in different regions of phase space we have determined the *effective* (i.e. in-medium) K-factor, defined as the ratio of the inclusive distribution for charm production in A-A collisions to the leading-order distribution without nuclear effects. We have shown that the K-factor, in general, is a function of rapidity, transverse momentum and x_F ($x_F = p_z/p_{zmax}$, and p_z is taken to be in the beam direction) and can not be taken to be a constant.

Our results, which include the next-to-leading order corrections, $O(\alpha_s^3)$, and the nuclear shadowing effect for the differential and total inclusive cross sections for charm production in hadronic and heavy-ion collisions at RHIC and LHC energies can be found in Ref. 6 and 7. We have shown that low-energy data for the total cross section for charm production in p-p and p-A collisions provide some constraints on out theoretical parameters, especially for the choice of the renormalization scale. The choice $Q = m_c$ seems to be preferred by the data. Theoretical uncertainty in the calculation of charm production in nuclear collisions due to the choice of the structure function is small at Fermilab fixed target energies (only few percent) and at RHIC (8% in the central rapidity region), while at LHC energies this uncertainty is about a factor of 6. Furthermore, the shape of the rapidity distribution for charm production in Au-Au collisions at LHC is very sensitive to the low-x behavior of the gluon structure function, resulting in a larger "dip" at $y = 0$ for a more singular function. Similarly, the transverse momentum distribution at LHC is much steeper when obtained with MRS D− structure function than with less singular structure function, such as MRS A set or the non-singular structure function, such as MRS D0 set. For the rapidity distribution, we have found the hadronic K-factor to be about 2 (2.5) at RHIC (LHC) energies, while the *effective* K-factor for Au-Au collisions is about 1.5 (1.1). In case of p_T distribution, the hadronic K-factor at LHC energies varies from 1.4 at $p_T = 0.7$ GeV to 7.6 at $p_T = 6$ GeV, while the *effective* K-factor changes from 0.6 to 4.6 in the same p_T range. This behavior of the *effective* K-factor is a direct consequence of the fact that the low p_T region (or small x) corresponds to the maximum shadowing of the gluon distribution, which in case of gold is about 62%, while the larger values of p_T probe region of phase space where the nuclear shadowing is smaller. We have also obtained the *effective* A-dependence for the differential

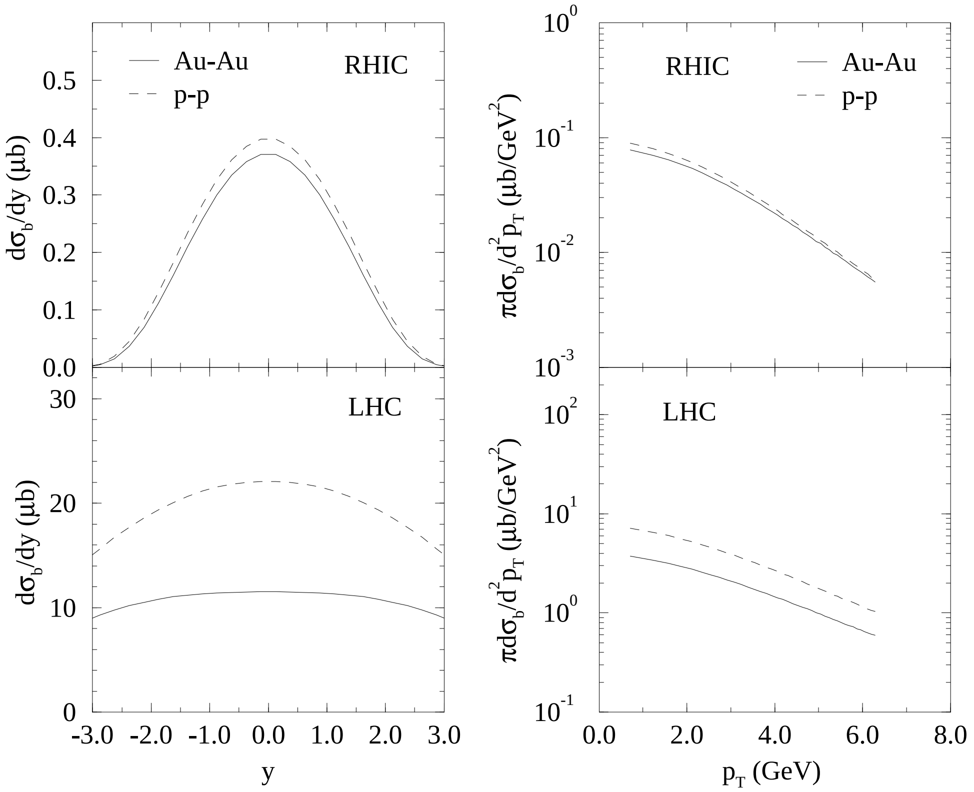

Figure 2.2: The rapidity and p_T distributions for bottom quark production in p-p and Au-Au Collisions at RHIC and LHC energies. The Au-Au cross sections are scaled by $1/A^2$ so a comparison to p-p collisions can be made. These distributions are calculated to $O(\alpha_s^3)$ with MRS A structure functions.

and total cross sections for charm production in p-Au collisions at RHIC and in Au-Au collisions at the LHC. We have found that the dominant contribution to charm production comes from the initial state gluons (about 95% at RHIC and 99% at LHC). Thus, combined measurements of inclusive charm production in hadronic and nuclear collisions at these energies, in addition to providing an important test of perturbative QCD in the small-Q^2 and small-x region, might be able to provide valuable information about the elusive role of gluons inside a nucleus, especially in the region of very small x.

Future hadronic and heavy-ion experiments at RHIC and LHC, PHENIX (RHIC), ALICE (LHC), and CMS (LHC), are currently planning to measure the open charm production. The PHENIX experiment will be looking for dilepton coincidences as a signature for charm production in the central rapidity region ($y \approx 0$) and at higher rapidity ($y \approx 2$). ALICE will have full acceptance in the central rapidity region, $|\eta| < .9$. The CMS is planning on full muon detection up to $|\eta| = 2.4$. Each of these groups has expressed interest in our work, and has asked us to provide information on the background processes. One of the signatures for charm production is di-muon coincidences. The primary backgrounds for this process is $B\bar{B} \rightarrow \mu^+\mu^-$, so it is important to compare the bottom production to the charm predictions. Since mass of the bottom quark is large compared to the QCD scale, we expect perturbative QCD calculations to be even more reliable than in case of the charm. We have calculated rapidity and p_T distributions for bottom quark production in p-p and Au-Au Collisions at RHIC and LHC energies. These distributions are displayed in Fig. 2.2.

In our calculations we have taken the bottom quark mass to be $m_b = 5GeV$. We find that at RHIC(LHC) energies the number of bottom quark pairs produced per unit rapidity in central Au-Au collision in the central rapidity region is 0.01 (0.3), while the total number is .03 (2.8). The *effective* (nuclear) K-factor for the total bottom cross section at RHIC (LHC) is 1.7 (1.4), reflecting significant shadowing at LHC energies only. In case of p_T distribution, we find that nuclear K-factor has strong p_T dependence only at LHC energies, ranging from 1.1 at $p_T = 0.7$ GeV to 1.4 at $p_T = 6$ GeV.

Since January 1995 (after Sarcevic's talk at Quark Matter '95), we have been approached by several members of the experimental heavy-ion collaborations (PHENIX, STAR, CMS, ATLAS, and ALICE) and we are providing them with necessary theoretical input for detection of heavy quarks.

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2.2 Gluon Density in a Nucleus at Small x

(I. Sarcevic, P. Valerio and I. Zubarev)

In the previous section, we have discussed the large theoretical uncertainties in the predictions for the heavy quark production at LHC energies due to the unknown low-x behavior of the gluon density in a nucleon and in a nucleus [1]. At these energies, the dominant sub-process for heavy quark production is gluon fusion. Currently there are no direct measurements of a gluon distribution in a nucleus. Furthermore, for the x-values probed at these energies ($x \sim O(10^{-4})$ at rapidity $y \approx 0$ and $x \sim O(10^{-6})$ at rapidity $y \approx 5$) there is also no experimental data on gluon structure function in a nucleon [2]. In our calculation of the heavy quark production in heavy-ion collisions, we have assumed that each parton structure function in a nucleus can be written as the product of a parton structure function in a nucleon and an appropriate function that models the nuclear effects. Since gluon density in a nucleus depends in a non-trivial way on the gluon density in a nucleon, better understanding of the low-x behavior of the gluon density in a hadron *and* in a nucleus is needed.

2.2.1 Parton Evolution Equations: Current Status

In this section, we give a brief review of the current status on evolution equations for the parton distributions. There are two main issues that we want to address. First, we will determine the correction to the anomalous dimension of the twist-$2n$ quasi-partonic operator in the Wilson operator product expansion (OPE) of the gluon density, due to the 3-pomeron direct interaction. This particular higher-twist correction is going to modify the existing parton evolution equations. Furthermore, this correction is expected to be important in the kinematic regions that will be accessible to HERA and LHC, and it has ramifications in both hadronic and nuclear structure function evolution at small-x. Secondly, we will use this new parton evolution equation to evaluate the gluon density in a nucleus by using appropriate initial conditions, as elaborated below. This calculation will provide insight into the experimentally measured nuclear shadowing effect [3].

For the calculation of perturbative QCD processes involving hadronic initial states, the parton model is the most successful approach. In this model, the non-perturbative effect of confinement is encapsulated into a set of quark and gluon structure functions, which depend on the kinematic variables x and Q^2, representing the momentum fraction of the hadron carried by the parton and the virtuality of the probing photon. The most simple parton model description of deep inelastic lepton-hadron scattering leads to the so-called Bjorken Scaling of parton distributions. That is, the structure functions are independent of Q^2 to first order. However, when higher order QCD corrections are taken into account, such as gluon bremsstrahlung from the initial parton states, the scaling behavior is violated, and the structure function depends logarithmically on Q^2 as well as on x. At large enough Q^2, the expansion parameter, $\alpha_s \sim (\ln Q^2)^{(-1)}$, is small enough that perturbative QCD (PQCD) is expected to be valid. However, when calculating the bremsstrahlung corrections to deep inelastic scattering, terms of the form $\alpha_s \ln Q^2$ appear, and therefore α_s is not an appropriate expansion parameter. In this case all diagrams which have $(\alpha_s \ln Q^2)^n$ contributions need to be resummed. In the physical (axial) gauge, it has been shown that the dominant contribution comes from ladder diagrams, with strongly ordered transverse momentum [4]. By summing up the parton ladder diagrams , the evolution of parton densities can be described in the large-Q^2 and intermediate-x region by the Gribov-Lipatov-Altarelli-Parisi equations[4, 5]:

$$\frac{dq_i(x,t)}{dt} = 4\pi\alpha_s(Q^2) \sum_{partons} \int_x^1 \frac{dy}{y} P_{ji}(x/y) q_j(x,t) \tag{2.1}$$

where $t = \frac{1}{2}\ln(Q^2/Q_o^2)$ and $P_{ji}(x/y)$ is the probability of finding a parton of type i with momentum fraction x inside a parton of type j with momentum fraction y.

Ladder diagrams with vector particles in the t-channel contain a large $\ln(1/x)$ contributions for each α_s. The resummation of the leading $\alpha_s \ln(1/x)$ terms corresponds to summing gluon ladders . This resummation provides the description of structure function evolution to small-x, through the Balitskii, Fadin, Kuraev and Lipatov equation (BFKL) [6]:

$$\frac{\partial\phi(x,q^2)}{\partial\ln(1/x)} = \frac{N_c\alpha_s}{\pi}\int K(q^2,q'^2)\phi(x,q'^2)\frac{d^2q'}{2\pi}, \tag{2.2}$$

$$\phi(x,q^2) = \frac{\partial xG(x,q^2)}{\partial q^2}, \tag{2.3}$$

and

$$K(q^2,q'^2)\phi(x,q'^2) = \frac{\phi(x,q'^2)}{(q-q')_t^2} - \frac{q'^2_t\phi(x,q^2)}{(q-q')_t^2(q'^2_t+(q-q')_t^2)}, \tag{2.4}$$

where N_c is the number of colors. The first term in the kernel, $K(q^2,q'^2)$, describes the emission of a gluon, and the second term describes the reggeization of all of the t-channel gluons in the ladder. This equation is valid only for small-Q^2 and small-x.

The solution to the BFKL equation increases rapidly at small-x (i.e. $xG(x,Q^2) \sim x^{-0.5}$ for small x). This behavior is in agreement, although it is on the outer limits of the new HERA data, which can be fitted with [7]:

$$F_2^{exp} \sim (\frac{1}{x})^{\omega_{exp}}, 0.3 < \omega_{exp} < 0.5. \tag{2.5}$$

Thus, the BFKL equation is a reasonably good starting point in the study of low-x physics. However, at small enough x, this singular behavior will result in the violation of the Froissart bound. This violation of unitarity, one of the most important properties of the S-matrix, can be seen in this way. The cross section for a virtual photon interacting with a nucleon must be less than the transverse size of the nucleon:

$$\sigma(\gamma^* N) < \pi R_N^2, \tag{2.6}$$

and this cross section can be written as

$$\sigma(\gamma^* N) = \sigma(\gamma^* g)xG(x,Q^2) \sim \frac{\alpha_s(Q^2)}{Q^2}xG(x,Q^2). \tag{2.7}$$

Taking a gluon density given by the solution to the BFKL equation, we can make an estimate of the kinematic region for which the BFKL solution will not violate unitarity:

$$\frac{\alpha_s(Q^2)}{Q^2}(\frac{1}{x})^{\omega} \leq \pi R_N^2 \tag{2.8}$$

From this equation we can extract a critical line in the $(\ln(1/x), \ln Q^2)$ plane, below which, these perturbative techniques are valid, and above which, these solutions violate unitarity:

$$\ln \frac{1}{x_{cr}} = \frac{1}{\omega} \ln \frac{\pi R_N^2 Q^2}{\alpha_s(Q^2)}. \tag{2.9}$$

For those x-values smaller than x_{cr} the BFKL approach needs to be modified.

The first attempt at modifying this equation was made over a decade ago by Gribov, Levin and Ryskin (GLR) [8]. They have taken into account the parton overlapping effect in dense systems. When partons overlap spatially, there is a chance for parton recombination, which effectively lowers the parton density. They have modified the BFKL equation by adding a non-linear term, corresponding to gluon recombination. This modified BFKL equation is called GLR equation and is given by

$$\frac{\partial \phi(x, q^2)}{\partial \ln(1/x)} = \frac{N_c \alpha_s}{\pi} \int K(q^2, q'^2)\{1 - \alpha_s \frac{\phi(x, q'^2)}{\phi_0}\}\phi(x, q'^2) \frac{d^2 q'}{2\pi}, \tag{2.10}$$

where $\phi_0 = \frac{32 N_c}{9 Q_0^2}$. In the Double Log Approximation (DLA), terms containing large $\alpha_s \ln Q^2 \ln(1/x)$ are summed. In the axial gauge, this corresponds to gluon ladders where the t-channel gluons have strongly ordered transverse momenta. The GLR equation also includes multi-ladder "fan" diagrams. In the DLA, the GLR equation can be written as:

$$\frac{\partial^2 x G(x, Q^2)}{\partial \ln \frac{1}{x} \partial \ln Q^2} = \frac{\alpha_s N_c}{\pi} x G(x, Q^2) - \frac{\alpha_s^2 \gamma}{Q^2 R_N^2} (x G(x, Q^2))^2, \tag{2.11}$$

where $\alpha_s \gamma$ is the three-pomeron vertex coupling, and $\gamma = \frac{81}{16}$ for 3 colors [9].

One of the assumptions of the GLR equation is that the gluons are uniformly distributed on a disc of size R_N, so then, when we define a parameter $\rho \equiv \frac{xG(x,Q^2)}{\pi R_N^2}$, the probability that two gluons have the same value of x and Q^2, denoted $P^{(2)}$, is approximately ρ^2, which appears in the second term on the right side of Eq. (2.11). This is valid only at large x and/or small Q^2. In general, $P^{(2)}$ and ρ^2 are related by [10]:

$$\frac{P^{(2)}}{\rho^2} \propto e^{\frac{1}{(N_c^2-1)^2} \sqrt{\frac{16 N_c \alpha_s}{\pi} \ln \frac{Q^2}{Q_0^2} \ln \frac{1}{x}}}, \tag{2.12}$$

where Q_0^2 is the initial virtuality of the cascade, on the order a few GeV. The GLR equation needs to be modified to include two-gluon correlations, $P^{(2)}$, instead of ρ^2.

At smaller values of x, multi-gluon correlations, $P^{(n)}$ becomes important. These multi-gluon correlations are related to the higher-twist terms in the Wilson operator product expansion [11]. To see this we write the gluon structure function using the OPE as:

$$x G(x, Q^2) = x G^{(1)}(x, Q^2) + \sum_n \frac{1}{(Q R_n)^{2(n-1)}} x^n G^{(n)}(x, Q^2). \tag{2.13}$$

The n-gluon correlation is related to the twist-$2n$ term in the OPE by $P^{(n)} = \frac{x^n G^{(n)}(x,Q^2)}{(\pi R_N^2)^n}$. In analogy to the GLR equation, evolution equations for the n-gluon correlations can be written as [12]:

$$\frac{\partial^2 P^{(n)}(x,Q^2)}{\partial \ln \frac{1}{x} \partial \ln Q^2} = \gamma_{2n}\omega P^{(n)}(x,Q^2) - n\frac{\alpha_s^2 \gamma \pi}{Q^2} P^{(n+1)}(x,Q^2), \tag{2.14}$$

where γ_{2n} is the anomalous dimension of the twist-$2n$ quasi-partonic operator, $\omega = N-1$, and N is the Mellin-conjugate variable to x. The anomalous dimension calculated in DLA approximation is given by [13]:

$$\gamma_{2n} = \frac{\overline{\alpha}_s n^2}{\omega}\{1 + \frac{\delta^2}{3}(n^2-1)\}, \tag{2.15}$$

where $\overline{\alpha}_s = \frac{\alpha_s N_c}{\pi}$ and $\delta = (N_c^2 - 1)^{-1}$.

Laenen and Levin [12] combined this infinite set of coupled differential equations into one differential equation. They start by defining a generating function:

$$g(x,Q^2,\eta) = \sum_{n=1}^{\infty} e^{n\eta} x^n G^{(n)}(x,Q^2). \tag{2.16}$$

The gluon density can be obtained from the generating function by:

$$xG(x,Q^2) = Q^2 R_N^2 g(x,Q^2,\eta = -\ln(Q^2 R_N^2)). \tag{2.17}$$

Then, the equations (2.14) can be written in the following form:

$$\frac{\partial^2 g(x,Q^2,\eta)}{\partial \ln \frac{1}{x} \partial \ln Q^2} = \overline{\alpha}_s \frac{\partial^2 g}{\partial \eta^2} + \frac{\overline{\alpha}_s \delta^2}{3}(\frac{\partial^4 g}{\partial \eta^4} - \frac{\partial^2 g}{\partial \eta^2}) - \alpha_s^2 \gamma e^{-\ln(Q^2 R_N^2)} e^{-\eta}(\frac{\partial g}{\partial \eta} - g). \tag{2.18}$$

This equation is a fourth-order PDE in 3 dimensions and it is difficult to solve. It is interesting to note that R_N is a parameter that can not exceed the size of the hadron, but it could also be related to the size of localized regions of high gluon density, so-called "hot spots". Since this equation takes into account multi-gluon correlations it is appropriate for dense nuclear systems, provided appropriate boundary and initial conditions are used.

2.2.2 Higher Twist Corrections to the Parton Evolution Equation

At the low-x, low-Q^2 kinematic region of our interest, the higher-twist corrections to the evolution equation may play a crucial role. We are currently calculating the 3-pomeron direct

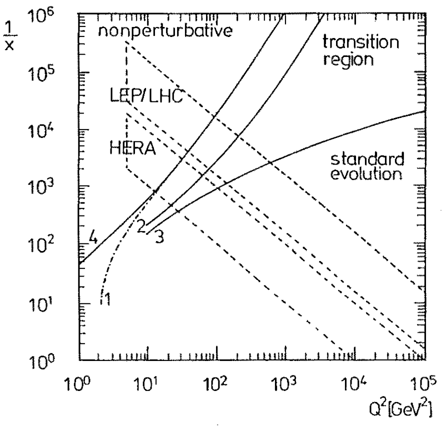

Figure 2.3: The kinematical regions expected to be probed at future HERA and LHC experiments. The curve "1" is the critical line given by (2.19).

interaction vertex and investigating its effect on the solution of the non-linear evolution equation given by Eq. (2.18).

We base our calculation on the effective pomeron field theory developed by Lipatov (see [14]), which allows us to obtain the anomalous dimension of the higher twist operators. In this approach, the pomeron interaction is described by an effective Hamiltonian, with a four-point pomeron interaction vertex and a relatively simple pomeron propagator. The anomalous dimension of the twist-$2n$ quasi-partonic operators, obtained for the effective pomeron field theory calculation, are essential in the GLR-like equations for the $P^{(n)}$. They play the role of the GLAP splitting functions. So far they have been calculated using only 2-pomeron correlations. Thus, the expression for γ_{2n} is exact only for $n = 1, 2$. At $n = 3$, there is a correction to γ_{2n} from the 3-pomeron correlations. The size of this corrections is presently unknown. In analogy to the analysis of [13], where the problem is cast in the form of a 2-dimensional Schrödinger equation, we plan to calculate the 3-pomeron direct interaction, using the same formalism. The kinematic region of the proposed HERA and LHC experiment will be partially in the non-perturbative region above, but near, the current critical line [15], as shown in Fig. 2.3. We believe that it will be important to get a handle on these corrections as the higher-twist effects begin to manifest themselves.

2.2.3 Solving New Evolution Equation for Large Nuclei

Once the higher-twist corrections are incorporated into the evolution equations, we can proceed to its solution. The initial boundary condition for the gluon density can be obtained by modeling the nucleus as a large collection of charmonium bound states. Using this formalism, Mueller showed that in the leading $\ln\frac{1}{x}$ approximation the nuclear structure functions can be written in a Glauber form [16]:

$$xG_A(x,Q^2) = \frac{16}{\pi^2}\int_x^1 \frac{dx'}{x'}\int\frac{d^2b_t}{\pi}\int_0^{Q^2} dQ'^2\{1 - e^{-\frac{3\alpha_s\pi^2}{4Q'^2}x'G_N(x',Q'^2)S(b_t^2)}\}, \tag{2.19}$$

where $S(b_t^2)$ is the nuclear density profile as a function of impact parameter. It is interesting to point out that this form is not factorizable i.e. $xG_A(x,Q^2)\neq R(x,Q^2,A)G_N(z,Q^2)$.

The analytical and numerical study of the new evolution equation (2.18), which incorporates higher-twist effects, will allow us to obtain the gluon distribution in low-x and low-Q region. Furthermore, we will be able to obtain the previously unknown A-dependence of the nuclear structure functions. We plan to start with the Glauber form (2.19) with new hadronic structure function, obtained as a solution of our new evolution equation. Then we need to evolve this nuclear gluon distribution using the same equation. Although this equation does not have an explicit A-dependence, it is non-linear so the A-dependence of the initial distributions will be translated in a non-trivial manner to the evolved structure functions. We expect the evolved gluon density to be non-factorizable. This will be the first theoretical calculation of nuclear gluon density that includes the correction due to the 3-pomeron interaction in the evolution equation. We expect that the non-trivial A and x dependence will also provide explanation for the observed shadowing effect of nuclear structure functions [3].

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2.3 Disoriented Chiral Condensate in Heavy-Ion Collisions

(Z. Huang, I. Sarcevic, M. Suzuki, R. Thews and X.-N. Wang)

The dynamical process of chiral phase transition is presently not well understood among many theoretical uncertainties in the description of the transition from the quark-gluon plasma to the hadronic gas. If there is significant supercooling before the plasma hadronizes, the evolution of the chiral order parameter is likely to be out of equilibrium. Recently it has been suggested that in ultrarelativistic heavy-ion collisions the rapid expansion of the collision debris in the longitudinal (beam) direction leads to a rapid cooling in the interior of the interaction region, and as a result, domains of the "unconventionally" oriented vacuum configurations allowed by the chiral symmetry may be formed [1]. Detection of this interesting phenomenon, the so-called Disoriented Chiral Condensate (DCC) [1], would provide valuable information on the vacuum structure of the strong interaction and the nature of the chiral phase transition.

The linear σ-model is shown to characterize the basic feature of the chiral phase transition since it is in the same universality class as QCD [2]. It also gives a good description of soft pion dynamics such as pion-pion scattering which dominates the late stages of evolution of the system. Preliminary theoretical investigations on nonequilibrium dynamics using the classical linear σ-model have found some evidence for the growth of long wavelength pion modes, which may indeed lead to domain structure or cluster formation [3]. Although the precise dimension of a typical DCC domain or cluster is currently under active debate, depending on the dynamical production mechanisms, it seems likely that in heavy-ion collisions many domains or clusters may coexist in the interaction volume.

The domain of misaligned condensate constitutes a macroscopic coherent pion filed with a certain orientation in isospin space. The DCC relaxes back to the true vacuum by radiating pion quanta with certain isospin correlation. It is predicted [1] that the probability distribution of neutral pion faction in total pions radiated from a DCC domain follows $P(f) = 1/2\sqrt{f}$, a very distinctive distribution compared to a binomial distribution. Other possible signals such as the dilepton emissions and multipion correlations have been suggested.

The search for a disoriented chiral condensate has been initiated by J. Bjorken and C. Taylor and their collaboration. A dedicated experiment, known as Minimax in Fermilab Tevatron collider, is well under way. Minimax strives to measure event-by-event the number of photons and number of charged particles traversing the acceptance. It is designed to explore the very forward region of phase space centered at a pseudorapidity of about 4 where anomalies of charge/neutral particle multiplicity distributions have been persistently reported by the cosmic-ray community. The first run is finished. About 2 million events have been recorded by the Minimax detector and the data analysis is being conducted.

Almost all experiment groups in the BNL RHIC collider including PHENIX, STAR, PHOBOS and BRAHMS plan to conduct the search for disoriented chiral condensate. The emphasis has been on the measurement on the p_t distribution of the pion multiplicity and two pion correlations which are sensitive to the existence of a DCC.

The topic of disoriented chiral condensate cannot live or die on theory alone; input from experiment is absolutely essential. We expect experimental results from Minimax sometimes next year after its run in winter 1996 with a lower beam energy. Interpretation of data is anticipated either for a discovery or for setting an experimental limit. This initial study done by Minimax is necessary to better understand experimental issues involved in DCC searches in heavy-ion collisions. In the meantime, there are many outstanding theoretical issues that need to be addressed.

2.3.1 Wavelet Analysis of DCC in Heavy-Ion Collisions

The numerical simulations of a DCC have suggested that the typical size of a DCC domain is not large, about 1-2 fm in transverse dimension and about 1 unit in spatial rapidity. In high energy heavy-ion collisions where the interaction volume is much larger than a DCC domain, many domains or clusters could be formed. If there are many uncorrelated small domains, the integrated probability distribution of the neutral pion fraction f emitted from a disoriented region, predicted to be $P(f) = 1/2\sqrt{f}$ [1], would become Gaussian. This follows from the Central Limit Theorem in statistics which states that the probability distribution of the sum of N independent identically distributed variables becomes Gaussian with width $\sim 1/N$ as N increases, independent of the original distribution. This makes the experimental search for the DCC signal a rather difficult task. In order to disentangle the DCC domain structure, we propose a new method which emphasizes not only the behavior of the probability distribution in the full phase space region but also its fluctuation in rapidity η or azimuthal angle ϕ. In other words, one needs to study the "local" properties of the distribution in phase space if the DCC clusters are "localized" objects in phase space.

Normally DCC domains are localized in coordinate space. If they develop collective motion in the course of their time evolution, they should also appear localized in momentum space. The standard "count in cell" technique is to divide the phase space into segments with a given acceptance size $\Delta\eta\Delta\phi$ and to define f *locally* in each cell. Since the location of a DCC cluster in phase space will not be fixed event by event, it is not useful to examine the probability distribution of f in one particular cell. Instead, one may look at the probability of finding a particular value of f in *all* cells of the same size in one event and average over the events. Clearly, the probability distribution will depend on the size of the cell. In addition, the information on the fluctuations with a characteristic correlation scale smaller than the acceptance is entirely suppressed.

We would like to propose a multiresolution analysis performed by a discrete wavelet

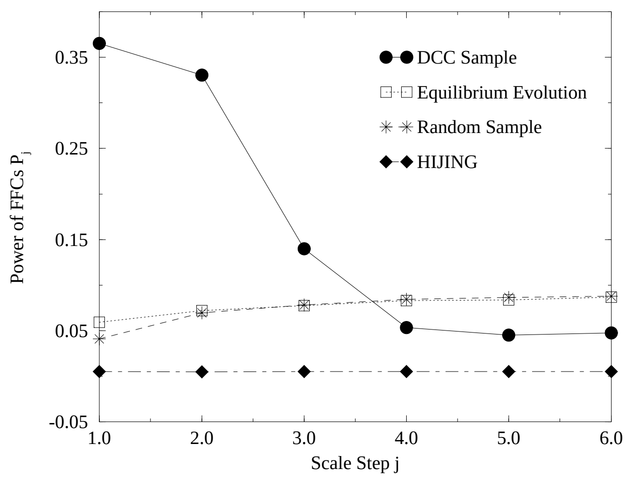

Figure 2.4: The power spectra of the wavelet coefficients for different dynamical scenarios.

transformation (DWT) which has been found effective in systematically detecting structures on various scales in turbulence, astrophysics, and multiparticle productions [4]. We demonstrate that the DWT proves to be very useful in identifying and measuring the DCC domain structures *simultaneously* in terms of their size (in scale) and location (in space). Since it is likely that there are other physical scales accompanying the typical DCC domain scale in a physical process, the multiresolution feature of the DWT is essential for identification of the structures of interest. It acts like a mathematical microscope which can zoom in or out to various scales at each location. Due to the completeness and orthogonality of the DWT basis, there will be no information loss.

We have conducted a preliminary study [5] using the wavelet method to analyze the fluctuation of the neutral pion fraction f in the spatial rapidity in numerical simulations of a classical linear σ-model in 1+1 dimensions [3]. The probability distributions of the neutral pion fraction for various rapidity-bin sizes have distinctive shapes in the case of a DCC and exhibit a delay in approaching the Gaussian distribution required by the Central Limit Theorem. We also find the wavelet power spectrum for a DCC, shown in Fig. 2.4, to exhibit a strong dependence on the scale while an equilibrium system and the standard dynamical models such as HIJING have a flat spectrum.

We would like to apply the wavelet method to analyze the experimental data when they become available. It seems promising that the DCC domain structure in heavy-ion collisions may be revealed by observing a *Delayed Central Limit* in probability distribution of wavelet transformed neutral pion fraction and a rise in wavelet power spectrum.

2.3.2 Studies of DCC Formation Mechanisms on Lattice

Theoretical investigation on the DCC formation mechanisms has been studied in the context of a linear and a nonlinear σ-models. It is generally believed that in the presence of explicit chiral symmetry breaking some nonequilibrium dynamics is needed in order for the disoriented region to grow into a sizable domain. Two dynamical scenarios have been suggested responsible for an unstable motion: quenching [2] and annealing [6]. It is shown by Akasawa, Huang and Wang [3] that in both cases the role of pion field fluctuations is very essential. In the quenching scenario, the instability is planted in the initial fluctuations of the pion field; while in the annealing case, the instability could only develop in the subsequent evolution of the system. Although some studies have favored a quenching over an annealing, no firm conclusion has been made.

Most of existent numerical work so far has simulated the initial fluctuations of pion field on the lattice using a Gaussian noise distribution with a momentum cutoff provided by the lattice size. However, such a white noise distribution is not *a priori* justified. We would like to propose to study a more realistic situation where the fluctuation is governed by a thermal distribution as it is more appropriate for a thermalized system expected to be produced in heavy-ion collisions. In contrast to a white noise fluctuation which has a flat momentum distribution, the thermal distribution has an enhanced low momentum distribution which may strengthen the instability as it arises from low momentum modes and thus lead to the DCC domain formation.

Another important aspect of the numerical simulation is the introduction of a minimal correlation length for the pion field through the introduction of the lattice spacing. This is a very nontrivial assumption which implies that the pion field is a constant vector within the lattice spacing in isospin space. Each lattice location is associated with a unique direction in isospin space and all possible directions are equally probable. We would like to investigate the sensitivity of the pion charge distribution on this assumption, in particular, whether the existence of large fluctuations in neutral pion fraction can be attributed to such a minimal correlation.

2.3.3 Multiparticle Correlations from a DCC Source

The pions radiated from a disoriented chiral condensate are coherent and have a very distinctive final charge distribution. Studying pion correlations in the case of DCC is equally interesting. We would like to investigate the impact of the DCC on the correlations of pion pairs with same and opposite charges. Bose-Einstein correlations of identical pions have been measured in many nuclear reactions [7]. Pions produced by conventional scattering mechanisms are largely incoherent, and exhibit an intensity interference that is analogous to the Hanbury-Brown-Twiss effect for photons. In contrast, a fully coherent source of pions exhibits interference at the amplitude level. Consequently, the HBT effect is to be suppressed in heavy-ion collisions if a large DCC domain forms. One may also study the correlations between the pions with opposite charges. The characteristic correlation over the azimuthal angle ϕ for pions having zero rapidity is given [1]

$$\frac{\sigma_{\rm tot}\frac{d\sigma^{\pi^+\pi^-}}{dk_1 dk_2} - \frac{9}{10}\frac{d\sigma^{\pi^+}}{dk_1}\frac{d\sigma^{\pi^-}}{dk_2}}{\frac{d\sigma^{\pi^+}}{dk_1}\frac{d\sigma^{\pi^-}}{dk_2}} = \frac{3}{10}\cos^2(\phi_1 - \phi_2). \tag{2.20}$$

The experimental verification of such predictions is very interesting. On the other hand, there has been no clear theoretical prediction on the correlations for pion pairs in the rapidity distributions except for an intercept in Ref. [8] derived from an isosinglet coherent state. We would like to study rapidity correlations in the case of a DCC from both a numerical simulation and an analytical method. Another interesting phenomenon is pion correlations originating from a DCC squeezed state [9]. It has been suggested that a quantum description of a DCC state from a quenching in the linear σ-model is a squeezed state. Depending on the squeezing parameter, the two-particle correlations can be very different from the conventional ones [9].

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2.4 Partial $U(1)_A$ Restoration in Heavy-Ion Collisions

(Z. Huang and X.-N. Wang)

At the Lagrangian level, QCD has, in addition to $SU(N_f) \times SU(N_f)$ chiral symmetry, an approximate $U(1)_A$ symmetry, under which all left-handed quark fields are rotated by a common phase while the right-handed quark fields are rotated by an opposite phase. It is well known that the $U(1)_A$ symmetry is violated by the axial anomaly present at the quantum level and thus cannot give rise to the Goldstone boson which would occur when $U(N_f) \times U(N_f)$ chiral symmetry is spontaneously broken. The $U(1)_A$ particle, known as $\eta'(958)$ in the $N_f = 3$ case, acquires an additional mass through the quantum tunneling effects mediated by instantons, breaking up the mass degeneracy with pions, kaons and η in the chiral limit when all quarks (u, d and s) are massless. The $\eta(547)$ particle also acquires an additional mass through the mixing with η'.

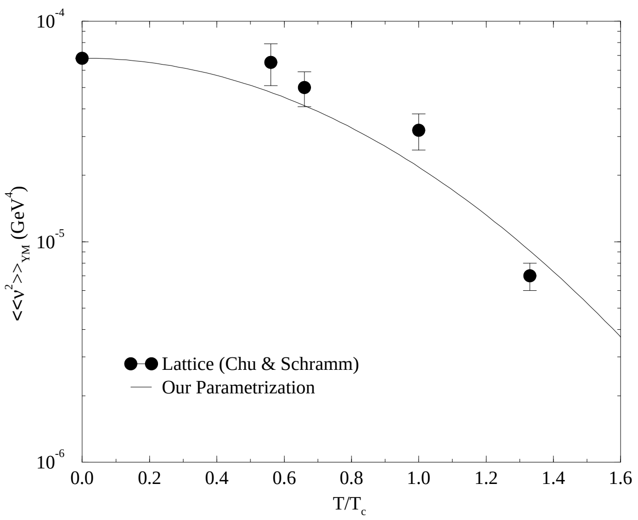

Figure 2.5: The lattice Calculation of the instanton density at finite temperature and a possible theoretical parameterization

It is believed that at high temperatures the instanton effects are suppressed due to the Debye-type screening [1]. Then one expects a practical restoration of $U(1)_A$ at high temperatures. The recent Lattice calculation [2] on the instanton density (the topological susceptibility $\langle\langle\nu^2\rangle\rangle$) in pure Yang-Mills theory at finite temperature has confirmed such an expectation. The lattice result is shown in Fig. 2.5. If the instanton effect becomes partially suppressed at a temperature lower than the chiral phase transition temperature T_χ, as indicated by the lattice calculation in Fig. 2.5, there may be some interesting phenomenological implications in high-energy heavy-ion collisions, as suggested first by Pisarski and Wilczek [3] and more recently by Shuryak [4]. One of the consequences of $U(1)_A$ restoration is the enhancement of η particle production at small and intermediate transverse momenta due to the softening of its mass at high temperatures. However, the final yield of the η particles and their p_t distributions both depend crucially on the chemical and thermal equilibrating processes involving the η.

Preliminary studies have been made by Huang and Wang [5] to evaluate the rates of various processes relevant for the thermal η particle production, in particular, whether or not the η can decouple early enough from the thermal system expected to be produced in relativistic heavy ion collisions. We have computed the thermal cross sections for the

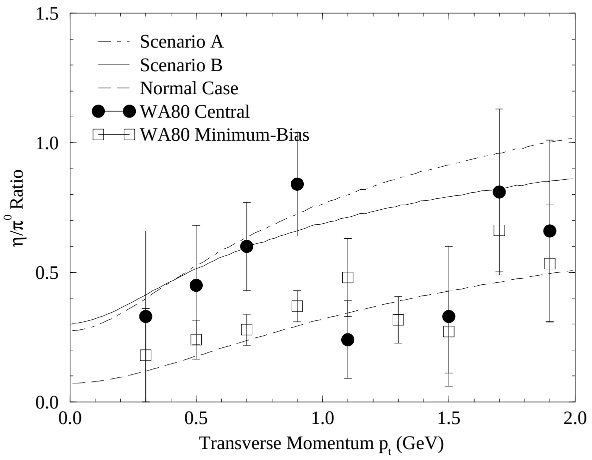

Figure 2.6: The predicted ratio η/π^0 as a function of the transverse momentum p_t in three scenarios as discussed in the text. The preliminary data from WA80 are also indicated.

processes $\eta\eta \leftrightarrow \eta\eta$, $\pi\eta \leftrightarrow \pi\eta$ and $\eta\eta \leftrightarrow \pi\pi$, essential to the thermal and chemical equilibration. The calculation is based on models which explicitly incorporate the $U(1)_A$ anomaly. An exponential suppression of the $U(1)_A$ anomaly due to the Debye-type screening of the instanton effect [1] is parameterized according to the lattice calculation shown in Fig. 2.5, which leads to the temperature dependence of the η and η' masses. It is found that the chemical equilibrium between η and π breaks up considerably earlier than the thermal equilibrium. Two distinct scenarios (scenarios A and B) for the η freeze-out are discussed. We have predicted a modest enhancement of thermal η production as a signal for the relic of $U(1)_A$ restoration.

The prediction for an enhanced η/π ratio is shown in Fig. 2.6 compared to the normal case where the zero-temperature η mass at the thermal freeze-out is assumed. The η/π ratio and its p_T dependence in the normal case is compatible with the data of minimum-biased $S + Au$ events from WA80 experiments where peripheral collisions dominate and we believe there is no partial $U(1)_A$ restoration. Comparing to this normal case, both Scenarios A and B predict an enhancement of the thermal η production due to a partial $U(1)_A$ symmetry restoration. Although preliminary data from WA80 on η/π^0 ratio in central $S+Au$ collisions at the CERN SPS energy, as indicated in Fig. 2.6, have shown such a trend of enhancement over the minimum-biased events, one certainly needs better statistics of the central collisions in order to make a definite conclusion.

As more central collision data for the η/π ratio will be available in future, there are many theoretical issues needed to be discussed. We would like to study in detail the elementary collision process for η and π and to reduce the theoretical uncertainties arising from the resonance poles in the scattering amplitudes. One possible way is to use the experimental information from hadronic collisions and to extract the relevant cross sections. A related topic is the effect of $U(1)_A$ symmetry restoration on η' particle production [6]. If the $U(1)_A$ anomaly effect is suppressed at high temperature, the mixing effect between η and η' is also suppressed such that η' is predominantly a $\bar{s}s$ state. Its abundance is therefore determined by the strangeness production in the dense hadronic matter.

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2.5 Thermal Dilepton Production in Heavy-Ion Collisions

(Z. Huang)

The lepton pair emissions are the ideal probes of the dense system of strongly interacting particles expected to be produced in ultrarelativistic heavy nucleus-nucleus collisions [1] Once produced, they will decouple from the strongly interacting system without further interactions and thus carry the information on the dynamical properties of the hot nuclear system. The theoretical predictions on the emission rate, both from a quark-gluon plasma at high temperature (T) and from a hot hadron gas at relatively lower T, have been extensively discussed in literatures [2, 3, 4, 5, 6]. There already exist some experimental data from various CERN SPS experiments (for a recent review, see [7]), and more thorough investigations will be made at upcoming RHIC and LHC experiments. The quantitative interpretation of the data is beginning to be possible [8, 9]

The main theoretical uncertainties in calculating the dilepton emissions rate come from the fact (apart from the crucial assumption that the system thermalizes in short time and stays in local equilibrium) that the system is strongly interacting and the perturbation theory fails. In a hot hadronic gas, though the effective degrees of freedom are mainly pions, pions interact strongly among themselves in many resonance channels, and effectively the resonances have to be taken into account. One obvious example is the pion electromagnetic form factor which exhibits the ρ-meson dominance. Near the resonances the naive pole approximations often violate the unitarity of the partial waves. There may also be some potential double counting problems associated with the emission rate at the location of the resonance peaks, as to whether or not the real resonance dilepton decays should be counted separately. Also, the multi-pion state emissions, e.g. $3\pi \to \ell^+\ell^-$ and $4\pi \to \ell^+\ell^-$ may not be so suppressed as they can proceed through the resonance poles, e.g. $3\pi \to \omega^*, \phi^* \to \ell^+\ell^-$ and $4\pi \to \pi a_1 \to \ell^+\ell^-$. Taking all these reactions into account by the kinetic theory is a difficult task though not impossible. More importantly, process involving hadrons in the final state such as $n\pi \to \ell^+\ell^- m\pi$ may have been entirely overlooked. Even in the phase of quark-gluon plasma where quarks and gluons are supposed to be weakly interacting, the high order corrections in α_s may not be so small.

The aim of proposed studies is to use and to develop the expression that allows the non-perturbative treatment of the strong interactions and avoids the detailed counting of the relevant reactions. Such an expression is known to exist [2, 3, 10] in the form of a low temperature expansion when T^2 is smaller compared to the dilepton invariant mass squared M^2. We would like to interpret the dilepton emission rate in terms of the spectral functions of some hadronic currents that are known experimentally in the low energy e^+e^- annihilation process and τ lepton decays. This approach is most relevant for the case of a hadronic gas. The recent development in calculating the correlation functions of vector currents in a low temperature pion gas using the PCAC and current algebra [11] enables us to obtain the important corrections to the rate with soft pions emitted together with the dilepton in the final state. Our main preliminary result for the dilepton emission rate, after the transverse momentum and the rapidity variable are integrated out, is [12]

$$\frac{dR_{\ell^+\ell^-}}{dM^2} = \frac{4\alpha^2}{2\pi} MTK_1(M/T)\left\{\rho^{\rm em}(M) - (\epsilon - \frac{\epsilon^2}{2})[\rho^{\rm V}(M) - \rho^{\rm A}(M)]\right\} \tag{2.21}$$

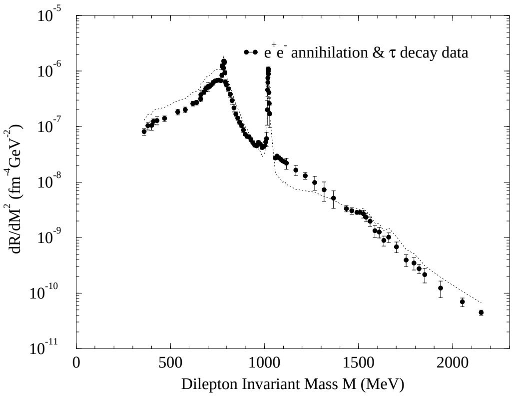

Figure 2.7: Dilepton emission rate from a $T = 160$ MeV thermal source. The dotted line (error bars not shown) represents the leading temperature contribution without the soft pion emissions. The error bars are calculated assuming the statistical independence of different experiment sets.

where M is the dilepton invariant mass, $\epsilon = T^2/6F_\pi^2$, T is the temperature and K_1 is the Bessel function. The spectral functions $\rho^{(i)}(s)$ $(i = \mathrm{em}, \mathrm{V}, \mathrm{A})$ can be extracted from the experimental data without any theoretical uncertainties:

$$\rho^{\mathrm{em}}(s) = \frac{s\sigma(e^+e^- \to \mathrm{had})}{16\pi^3\alpha^2} = \frac{R(s)}{12\pi^2}\;; \tag{2.22}$$

$$\rho^{\mathrm{V}}(s) = \frac{s}{16\pi^3\alpha^2}\sum_{n=1}\sigma(e^+e^- \to 2n\pi)\;; \tag{2.23}$$

and

$$\rho^{\mathrm{A}}(s) = \frac{8\pi m_\tau^3}{G_F^2\cos^2\theta_c(m_\tau^2+2s)(m_\tau^2-s)^2}\sum_{n=1}\frac{d\Gamma(\tau\to\nu_\tau(2n+1)\pi)}{ds}\,, \tag{2.24}$$

where s is understood to be the invariant mass squared of $(2n+1)\pi$'s, G_F is the Fermi constant and θ_c is the Cabibbo angle. Our parameterization clearly exhibits the mixing phenomenon of vector and axial vector currents at finite temperature and its relation to the dilepton production. We have obtained a new compilation of the experimental data for the spectral functions in the energy range from 360 MeV to 2150 MeV, relevant to the dilepton yield from a hadron gas. The dilepton emission rate for a $T = 160$ MeV is plotted in Fig. 2.7.

To compare the theoretical prediction on the thermal dilepton production with experimental data, one has to integrate the emission rate over the space-time history of the thermal system where the temperature is function of time. CERN SPS has observed some enhanced dilepton yield with the invariant mass below 500 MeV [7]. It is very important to use our formulation to calculate the final dilepton spectrum as it includes many previously unaccounted reactions and to investigate how much the difference there is. We are currently working on the hydrodynamical models and expecting some results soon.

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2.6 QCD-based Effective Theory of Particle Production in Heavy-Ion Collisions

(I. Sarcevic, T. Elze and H. Eggers)

In the past year we have continued our work on developing a dynamical model for describing the unusually large multiparticle correlations, the so-called intermittency phenomenon, recently observed by several experimental groups [1]. Previously, we have shown that the low-energy hadronic data on factorial moments measured in different rapidity bin sizes can be described with conventional exponential two-particle correlations and the linked-pair *ansatz* for higher-order correlations, without invoking the assumption of intermittency or power-law behavior of the moments [2]. The "linked pair" structure of higher order correlations that we introduced, continues to be confirmed even in hadronic collisions at higher energies, although as yet there is no basic understanding of the origin of this simple structure [3].

Three years ago, the multidimensional analysis of the factorial moments in different phase space regions has been done by several experimental groups, showing no saturation effect in the small regions of phase space [1]. This interesting result has inspired us to continue theoretical study of this problem. By looking at statistically independent correlations, given in terms of cumulants, we have found that in multiparticle production in heavy-ion collisions there are *only* two-particle correlations present (no true higher-order correlations) [4], perhaps due to a strong collective phenomenon present in multihadron production. A similar conclusion was reached by Baym, Friedman and Sarcevic several years ago when NA34 transverse energy distribution in heavy-ion collisions was analyzed in the independent-collision model.

Recently, we have developed a three-dimensional statistical field theory of particle production in heavy-ion collisions [5], based on an analogy with the Ginzburg-Landau theory of superconductivity. Taking the effective action to be Gaussian, and assuming the "field" to represent multiparticle density fluctuation, results in a two-particle correlation function of Yukawa form and no higher-order correlations. We have found a general scheme for projecting the theoretical three-dimensional correlation functions onto data of three and fewer dimensions and illustrated how theoretical predictions are tested against experimental moments in the different dimensions [6]. We gave examples in terms of NA35 parameters which should be testable against future data.

Currently we are pursuing the connection between a Ginzburg-Landau-type model and a more fundamental field theory. Assuming that the observed multiparticle correlations and fluctuations arise in the dense hadronic phase of matter (after a possible hadronization phase transition), the first step of a derivation for this conservative scenario consists in choosing a phenomenologically satisfactory effective field theory of hadronic interactions such as the sigma model. (This presents, of course, an alternative to the view that the relevant fluctuations stem from partonic showering processes which might seem more natural for high-energy e^+e^--collisions.) Keeping the high excitation energy of the compressed matter in mind, one tentatively begins with a field theory at finite temperature, neglecting for simplicity the difficult non-equilibrium aspects of heavy-ion collisions. Aiming at a three-dimensional Ginzburg-Landau statistical picture one is led to integrate out the (imaginary) time-dependent modes from the path-integral representation of the (sigma model) field theory. So far, we have achieved this only in a one-loop approximation for an arbitrary scalar field theory yielding a temperature-dependent three-dimensional effective action. It represents the infrared limit of the originally chosen theory in accordance with ideas on dimensional reduction and can be used as a starting point to calculate the correlation functions studied in the present paper. Clearly, much more can be done here to include non-perturbative effects if one wants to investigate, for example, effects of a phase transition possibly contained in the chosen field theory. Formally integrating out the fields of the three-dimensional effective action in favor of the fields squared ("densities"), one obtains a Ginzburg-Landau "free energy" functional with its parameters, given as temperature-dependent functions of the renormalized coupling constants of the original field theory.

Having outlined these further steps in the development of our model, we conclude that the study of correlations and fluctuations in dense hadronic matter seems a promising approach to further our understanding of strong interactions. In particular, the investigation of one-particle observables alone, which are nicely reproduced by a wide selection of event generators, in principle cannot provide a sufficiently detailed knowledge of the most interesting (and least understood) soft physics aspects of QCD.

Recently, Sarcevic has been invited by World Scientific to write a review article on "Intermittency in High-Energy Collisions: Recent Developments", which she expects to complete in Spring 1996.

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2.7 Wavelet Correlations in Multiparticle Production

(P. Carruthers, M. Greiner, P. Lipa)

Multiparticle production is often dubbed in terms of branching processes. In this context, selfsimilarity and scaling are particularly appealing concepts since they constitute an intriguingly simple and universal organization principle of the underlying dynamics. They are found in a large variety of complex phenomena.

Simple discrete hierarchical branching models with selfsimilar density fluctuations were first introduced in the description of energy dissipation in fully developed turbulence [1]-[2]. They combine an independent hierarchical evolution of different branches with a random multiplicative nature of density fluctuations. In multiparticle processes such selfsimilar cascade models serve as paradigmatic toy models for intermittency [3]-[5]. Intermittency is now experimentally well established in soft multiparticle interactions and is now being explored in hard (perturbative) multiparton interactions.

Although these branching models describe local density fluctuations, they are conventionally analyzed (globally) in integrated form as is the case for the experimental data. The underlying concepts of scaling and selfsimilar fluctuations are revealed in a multifractal (moment) analysis [6], which globally averages over the differential (local) correlation structure. A fully differential correlation analysis itself would be a more fundamental approach to reveal selfsimilarity and scaling and deviations from these. However, correlation functions of higher than second order are difficult to measure and to visualize. Nevertheless, we prefer to stick to a fully local description of hierarchical branching models and their analysis in terms of differential correlation densities.

The call of wavelets [7]-[8] [15] needs some further introductory remarks: The whole wavelet basis is constructed from dilations and translations of one single "mother" wavelet and represents a selfsimilar and orthogonal basis. Thus we would expect that the wavelet transform uniquely simplifies the correlation functions of selfsimilar processes and, in particular, quasi-diagonalizes the covariance matrix. This has been demonstrated in ref. [9] for fractional Brownian motion and in [10] for the p-model which is the simplest tractable random cascade model [2]. Still more can be learned from the wavelet transform.

Wavelet correlations provide statistical information about further subclustering of certain structures like clusters, filaments or voids. Correlations between small subclusters living inside larger clusters or voids are naturally revealed in the wavelet transformed correlations, as are correlations between voids and filaments. Loosely speaking, we denote wavelet correlations as "clump" correlations.

We have studied a class of discrete hierarchical cascade models with what we call random or deterministic branching and restrict ourselves to the one-dimensional case only [11]. From a general evolution equation for the generating function we calculate recursion relations for conventional correlation densities. As specific examples, the correlation structure of the frequently used p-model [2, 12] and α-model [13, 3, 4] as well as the p-model with random branching and the QCD-motivated cascade model of ref. [16, 17] are analysed in more detail. We have also studied the importance of the conceptual distinction between moments referring to the forward evolution of the cascade and moments obtained from a backward analysis.

We would like to propose a refined study of the wavelet techniques in order for that to become a powerful tool for experimental analysis. In this respect wavelet packets might help to simplify the correlation structure even further. Also, in view of the recent success of correlation integrals, one should develop an analogue, so-called wavelet correlation integrals; this could set new standards of inferring information on clump correlations. Mathematically, the challenging problem is to define the wavelet transformation on spheres which might have important applications in the context of angular intermittency.

We believe that a true understanding of the intermittency phenomenon in multiparticle cascading is still a long way off. For general particle cascades it is not a priori clear that they follow a hierarchical evolution structure. Particle cascades should be formulated as point processes, which need not be hierarchically organized from the beginning. The relevant splitting functions should determine whether hierarchical structures could evolve dynamically. This might be the case for perturbative QCD. Work in this direction, including the use of the wavelet transformation, is presently being carried out.

More analytical studies of more realistic parton cascades within perturbative QCD are necessary in order to fully appreciate the outstanding features of wavelet correlations. The question "what is the best wavelet to choose?" for the description of the cascade process might be accessible with a transformation to forward cascade evolution equations instead of backward evolution equations, which have been employed so far. For the analysis of experimental multiparticle data more advanced wavelet techniques like a wavelet correlation integral have to be developed. A real jewel might become to determine wavelet correlations from Monte Carlo simulations of multiparticle processes and to switch on/off the hadronisation process; in this respect wavelets represent an ideal tool to investigate directly the heavily debated hadronisation process as the latter might only show up at the small scales of the wavelet correlations.

Apart from multiparticle physics, the above outlined wavelet technology might have applications in other fields of physics. For example, wavelet (clump) correlations may shed further light on the long-standing question in intermittent fully developed turbulence of how eddies decay at small scales and would allow to gain new and supplementary insight for phenomenological modeling. In disordered solid state systems multifractality and localization of electron wavefunctions are discussed; a wavelet analysis appears to be suitable. Last not least, the simplifying aspect of wavelet correlations might ignite the century-old inverse problem of how to extract dynamics from correlations with new fuel.

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2.8 Peripheral Mesons and Rapidity Gaps

(H. Lu and I. Sarcevic)

Large rapidity-gap jet events in deep-inelastic scattering have been observed at HERA *ep* collider. These events are characterized by the presence of a large region of the detector in the forward proton's direction with no hadronic particles. They have been interpreted as the scattering of a virtual photon off a hypothetical object called pomeron. The diffractive structure function based on this picture has been recently been measured at HERA [1].

Pomeron is usually treated as a virtual hadron emitted from the proton. The incoming virtual photon probes into the parton structure of the pomeron. Theoretical estimates

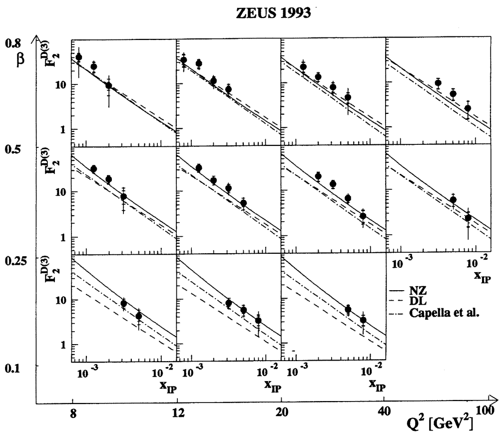

Figure 2.8: The experimental results of the diffractive structure function $F_2^D(\beta, Q^2, x_{I\!P})$ as measured by the ZEUS collaboration at HERA [1]. The inner error bars show the statistical errors, the outer bars correspond to the statistical and total systematic errors added in quadruture. Notice that in the low-$x_{I\!P}$ region, the distribution function follows a $1/x_{I\!P}^a$ behavior, represented by straight lines in the log-log plot.

for the pomeron content of the proton is far from certain. Ingelman and Schlein [2] have parametrized the pomeron content of the proton from a fit to hadron scattering data and obtained

$$f_{I\!P}(x_{I\!P}, t) = \frac{1}{4.6 x_{I\!P}} \cdot (6.38 e^{8t} + 0.424 e^{3t}), \tag{2.25}$$

where $x_{I\!P}$ is the momentum fraction of the proton carried by the pomeron, and t the momentum transfer squared. Donnachie and Landshoff [3] obtained another parametrization within the framework of Regge theory

$$f_{I\!P}(x_{I\!P}, t) = \frac{9\beta_0^2}{4\pi^2} F_1(t)^2 x_{I\!P}^{1-2\alpha(t)}, \tag{2.26}$$

with $F_1(t) = \frac{4M^2 - 2.8t}{4M^2 - t}(\frac{1}{1-t/0.7})^2$ being the proton's elastic form factor, $\beta_0 \simeq 1.8$ GeV^{-1} the pomeron-quark coupling, and $\alpha(t) = 1.085 + 0.25t$ the pomeron trajectory in Regge theory.

Deep-inelastic structure of the pomeron is also a highly uncertain quantity. Often simple parametrizations are chosen, like $F_2^{I\!P}(\beta, Q^2) = \frac{5}{3}\beta(1-\beta)$ or $F_2^{I\!P}(\beta, Q^2) = \frac{5}{3}(1-\beta)^5$ as used in the POMPYT Monte Carlo implementation [4] of the Ingelman-Schlein model, or like

$F_2^{I\!P}(\beta, Q^2) = 0.06\beta^{0.6}(1-\beta)^{0.6} + 0.015\beta^{-0.22}(1-\beta)^{4.6}$, as obtained by Capella *et al.* [5] by using Regge theory in the analysis of the diffractive structure function. Here, β is the Feynman variable representing the momentum fraction of the struck parton inside the pomeron, and Q^2 is the virtuality of the photon. We explore the interpretation of the large rapidity-gap jet events by using the meson-cloud picture of conventional nuclear physics. We study the peripheral meson degrees of freedom of proton, including the η, π, ω and ρ contributions. That is, the proton is interpreted as being surrounded by a meson cloud, with the core baryon being a nucleon (p, n) or an isobar $(\Delta^0, \Delta^+, \Delta^{++})$. We interpret the observed large rapidity gap events as coming from the deep inelastic scattering of the electron off the peripheral mesons. When the characteristic distance of the meson from the core baryon is larger than $\sim 1/\Lambda_{\rm QCD}$, the color coherence of the meson is separated from the color coherence of the core baryon. This would explain the presence of the large rapidity gap, since the partons of the peripheral meson hadronize far away from the remnant core baryon. In particular, by using the η, π, ω and ρ degrees of freedom, we can estimate the distribution of forward p, n and Δs in the large rapidity gap events. The observed $1/x^a$ behavior of the "Pomeron" structure function inside proton near $x = 0$, with a near unity: $a = 1.30 \pm 0.08(stat)^{+0.08}_{-0.14}(sys)$ (see Fig. 2.8), is interpreted as the dominance of peripheral vector mesons (ω, ρ) at small x, since the vector meson distribution in proton has a $1/x$ behavior at low x, while the pseudo-scalar mesons have a linear dependence on x. The deviation of a from unity is attributed to higher-mass excitations, like a f_2 meson surrounding an N^* core baryon.

As opposed to the case of the pomeron picture, in the peripheral-meson picture we are able to relate the rapidity-gap event cross section to traditional nuclear physics quantities, without assumptions on Regge theory or pomeron-quark coupling. Also, the deep-inelastic structure of pion serves as guidance to the structure function of other mesons. That is, in this approach, we have better control over the theoretical inputs, and much less assumption is needed in order to explain the observed rapidity gap events. Additionally, this approach will allow the estimate of the distribution of nucleons and isobars for the diffracted baryons.

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2.9 Particle Production in Sgr A* and in Active Galactic Nuclei (AGN)

(S. Markoff, F. Melia and I. Sarcevic)

One of the most exciting areas of current astrophysical investigation is the study of the interaction between matter and radiation in the extreme physical conditions encountered near massive black holes, particularly those in the nuclei of active galaxies. Several space agencies, including NASA and the European and Japanese Space Agencies, are undertaking missions designed to explore the X- and gamma-ray spectral characteristics of these sources. The expectation is that these observations will provide a unique probe of the environment close to the event horizon. The recently acquired data from current missions are finally enabling us to compare the predictions of theoretical models about such objects with their actual spectra. Using these data, we will construct a viable model for the high-energy radiative processes occurring near black holes in active galactic nuclei.

Currently we are looking at a specific black hole candidate, Sgr A* at the center of our galaxy, with plans to subsequently extend our results to its cousins in the nuclei of nearby galaxies, such as M31 and M104. The emphasis of the research is on the high-energy spectral component. This involves a thorough analysis of the micro-physical radiative processes and the interaction of matter and radiation under conditions impossible to recreate in any laboratory setting. A primary goal of our investigation is to anchor our theoretical calculations with recently acquired data from several international missions, including ROSAT, GRANAT, ASCA and the Compton GRO. Close association with the Infrared (IR) observers here at the University of Arizona will provide us with immediate ground-based IR observations of the galactic center. Additionally, our collaboration with members of the Whipple Observatory here in Tucson will give us access to observations of sources at TeV energies.

In the past few years, there have been several experimental developments regarding particle interactions that are crucial to our understanding of these high-energy processes. The current standard model assumes that hot, ionized gas from nearby stars is accreting onto the central engine via a Bondi-Hoyle bow shock formation. The gas heats up as it compresses toward the event horizon. A small fraction of the protons in this gas are then accelerated by the well-known Fermi process, in which particles are energized as they bounce back and forth across the shock interface, resulting in a power-law distribution of relativistic nuclei. Although the initial acceleration process is straightforward, the particle cascades which ensue from the interactions of the relativistic protons with the ambient matter and radiation require a more sophisticated model. Calculating these cascades accurately is crucial for correctly interpretating the radiation produced by the source.

After determining the complete particle spectra, taking into account all the important energy loss mechanisms, such as bremsstrahlung and cyclotron-synchrotron emission, we will calculate the radiative spectrum arising from all the relevant photon producing processes. These include: electron-electron and electron-proton bremsstralung, collisionless scattering of particles when the temperature of the electrons is the same as the ions and possibly when they are different, magnetic bremsstrahlung, high energy photons produced via pion decay in the particle cascades within the inflowing gas, and Compton scattering processes. Each of these produces a unique spectral signature. Knowing the cascades well and their contributions to these various processes, we expect that a comparison of the results from our calculations with the data will test our ideas about how black holes accrete, and whether they are consistent with our understanding of general relativity and the overall evolution of the nuclear regions of galaxies.

Currently we are modeling Sgr A*, which is the logical starting point for this project: it is nearby, and relatively easy to observe at IR, radio and gamma-ray wavelengths. The IR displays variability, perhaps suggesting a disk orbiting the black hole. Sgr A* is ideal for the initial stage of this project because it is a relatively weak source, which permits us to model its high-energy emission without having to take into account such complicating aspects as Compton scattering effects and radiative pressure. Thus we can test the basic premises of our model against the data, which we have in the form of an EGRET spectrum (John Mattox, private communication), without having to consider the possibility that radiation will have a feedback effect on the system dynamics.

The next stage involves using 3-D hydrodynamical simulations. From these we will obtain the particle number density, the temperature(s), the velocity and magnetic field intensity as a function of position within the accreting gas, and use these physical characteristics together with the emission processes described above to calculate the emissivity, radiative intensity, overall luminosity, and the time variability of Sgr A*'s spectrum as functions of the frequency.

After there is satisfactory agreement between the theoretical model and observations of Sgr A*, it will be essential to extend this investigation to the galactic nuclei, and thereby learn more about this class of "local" massive black holes, and how they relate to the more extended population of active galactic nuclei.

2.10 High Energy Neutrino Interactions of Relevance to Detection of AGNs

(I. Sarcevic, R. Gandhi, C. Quigg and M. H. Reno)

As the number of extant and planned experimental facilities for extra-terrestrial neutrino detection testifies, their importance in extending the particle physics frontier beyond the standard model is widely recognized. At the highest energies, neutrinos are decay products of pions produced in cosmic accelerators and thus provide a direct window to the most energetic processes in the universe. Neutrino astronomy also has important advantages over gamma-ray astronomy. In the energy range 10^{12} to 10^{20} eV, the expected air-cascade signature of the gamma-ray flux is dwarfed by the cosmic-ray background. In the same energy range the flux of neutrino by-products of the cosmic-ray interactions is very small, and the expected flux from cosmic accelerators should dominate. In addition, neutrino telescopes span a significant fraction of the sky at all times, in contrast to gamma-ray detectors which provide small angular coverage for a limited amount of time.

In general, fluxes of ultrahigh energy (UHE) neutrinos fall steeply with increasing energy, making the direct detection of these particles difficult. Cerenkov detectors, however, are capable of detecting upward-moving muons produced by the charged-current (CC) interactions of energetic neutrinos in the rock below the detector. The effective volume of the detector is thus enhanced in proportion to the range of the produced muon (typically several kilometers), which then may traverse the fiducial volume of the detector or stop therein.

The detection of UHE neutrinos via muons relies on the charged current interaction $\nu_\mu + N \to \mu + X$, where N is a nucleon. Due to the sharply falling neutrino flux with increasing energy, the usual mode of detection relies on increasing the effective detector volume by looking for muons produced in the rock below the detector. The expected event rate for these detectors is presently subject to two sources of uncertainty: *a*) the uncertainty in the knowledge of the incident neutrino flux, and, *b*) the calculation of the charged current cross-section for the $\nu - N$ interaction at ultra-high energies. The latter stems from necessary extrapolations of the nucleon quark structure functions for very low parton fractional momentum (x) and large momentum transfer (Q^2).

Using recent results from the $e-p$ collider at HERA [1] on small-x parton distributions and the improved theoretical knowledge on next-to-next-to-leading order corrections to the parton evolution to large Q^2, we have performed calculation of the neutrino-nucleon cross sections for neutrino energies up to $10^{20}eV$ [2]. For 10^{20}eV neutrinos we have found the cross sections to be factor of four to ten times larger than previous estimates. Since UHE neutrinos are detected via neutrino-induced muons, our results translate into significantly larger probability that neutrino-induced muon created in the rock below the detector would reach the detector and thus larger downward muon rates. For upward-moving muons, the neutrino attenuation in the Earth depends on the charged-current and neutral-current cross sections. For larger neutrino-nucleon cross sections, the attenuation effect is stronger. Since it is the product of the muon probability and the neutrino attenuation that goes into the calculation of the upward-moving muon event rate, one has to incorrporate these two competing effects, as elaborated below.

The generic mode of ultrahigh energy neutrino production in a cosmic point source is $p + matter \rightarrow \pi^\circ, \pi^\pm \rightarrow \mu^\pm, e^\pm, \nu, \bar{\nu}$. Candidate accelerators for the highest energy neutrinos include pulsars, blackholes and cosmic strings. Although cosmic rays of energies exceeding 10^{20} eV have been observed, their origin for energies above 10^{15} eV or so is not well understood, and neither is it presently possible to pin down the identity of a highly energetic primary. The primary particle in the recent Fly's Eye observation [3] of the highest energy cosmic ray ($\sim 3.2 \times 10^{20}$ eV) could be a proton or a neutrino with equal probability.

As is well known, the cosmic ray spectrum has two kinks, referred to as the "knee" (at around 10^{15} eV) and the "ankle" (at around 10^{19} eV). Below the knee, it is believed that the spectrum originates in shock wave acceleration of matter in galactic supernova remnants. Above these energies the spectrum is thought to originate in accelerators in active galactic nuclei (AGNs), both in our galaxy and extra-galactic.

Using the observed spectrum of UHE cosmic rays, the luminosity of the galaxy in protons of energy in the 10^{15} eV range can be estimated [4]. For an average source situated at roughly 10 kpc, one obtains a consequent neutrino flux

$$\frac{dN_\nu}{dE_\nu} = \frac{2 \times 10^{-9}}{E_\nu^2} cm^2 s^{-1} TeV^{-1} \tag{2.27}$$

where E_ν is in TeV.

Cosmic rays above 10^{18} eV cannot be accounted for by galactic accelerators. Nearby AGNs (~ 100 Mpc) could be the most likely origin of such UHE particles. Flux estimates from UHE point sources, especially extra-galactic ones, are a topic of active research at the present time [4, 5], with several unanswered questions. Question that we would like to ask would be the following: Given an incoming UHE neutrino of energy E_ν and a detector with a threshold for muon detection of E_μ^{th}, what is the probability $P(E_\nu, E_\mu^{th})$ that it produces a muon in the detector? Secondly, given a UHE neutrino flux from a particular source, what is the corresponding upward muon flux (including neutrino absorption in the earth)? We address these questions in what follows, with the viewpoint of highlighting the dependence of the answers on the CC cross-sections evaluated here.

Clearly, $P(E_\nu, E_\mu^{th})$ depends on the charged-current neutrino cross-section and the range of the muon in the earth (which will set the effective volume of the detector). The latter quantity is limited by the energy losses of muons in rock, due to ionization, pair and photo-production and bremstrahlung. The ionization losses depend only logarithmically on the energy, but the radiation losses are roughly proportional to the energy. For our purpose here, one may work in the "average loss" approximation in order to calculate the muon range. Parametrizing

$$-\,dE/dX = \alpha + \beta E \tag{2.28}$$

one uses the average values $\alpha \approx 2$ MeV g$^{-1}cm^2$ and $\beta \approx 3.9 \times 10^{-6}$ g^{-1} cm^2, with X being the distance (g/cm^2) traversed in the rock, to evaluate $P(E_\nu, E_\mu^{th})$. The consequent muon flux (ϕ_μ), incorporating neutrino absorption in the earth, and in terms of a given neutrino incident flux (ϕ_ν) from any source is

$$\phi_\mu(\Theta) = \int P(E_\nu) S(E_\nu, \Theta) \phi_\nu(E_\nu, \Theta). \tag{2.29}$$

Here the integral is over the incoming neutrino spectrum and Θ, as before, is the zenith angle.

The inclusive cross section for $\nu_\mu + N \to \mu^- + X$ is given by

$$\frac{d^2\sigma}{dxdy} = \frac{2G_F^2 M E_\nu}{\pi} \frac{M_W^4}{\left(Q^2 + M_W^2\right)^2} (xq(x,Q^2) + x(1-y)^2 \bar{q}(x,Q^2)), \tag{2.30}$$

where $x = Q^2/2M\nu$, $y = \nu/E_\nu$, with Q^2 being the momentum transfer between the neutrino and muon, and ν the energy loss in the lab frame, $\nu = E_\nu - E_\mu$. M is the mass of the nucleon and M_W is the mass of the W-boson, while the Fermi constant, $G_F = 1.16 \times 10^{-5}$ GeV^{-2}. With the assumption that the target is an isoscalar nucleon, the quark distribution functions contain valence and sea quark contributions.

At low energies, the neutrino-nucleon cross section reduces to the usual 4-Fermi approximation and increases linearly with the neutrino energy, with little dependence on the momentum transfer Q. At high energies, $E_\nu \gg M_W^2/2M$, the prediction for the total cross section in the 4-Fermi approximation violates unitarity. Instead, the W-propagator in Eq. (2.30) becomes important [6]. The resulting cross section depends on the momentum transfer, Q which appears in the quark distribution functions of a nucleon and in the W propagator. With larger Q, the parton distributions are enhanced, but the W propagator limits the growth of Q. For $E_\nu \gtrsim 10^5$ GeV, the average Q^2 is of the order of M_W, and the neutrino-nucleon total cross section is therefore dominated by the behavior of the quark densities at $x \sim M_W^2/2ME_\nu$ and $Q^2 \sim M_W^2$. For neutrino energies above 10^5 GeV, the small-x (i.e., $x \leq 3 \times 10^{-2}$) behavior of the quark densities plays an important role in determining the cross-section in Eq. (2.30).

Recently, knowledge of the parton densities at small values of x has been substantially improved by the striking observation of the dramatic increase of the proton structure function, $F_2(x, Q^2)$ with decreasing x for $10^{-4} \leq x \leq 10^{-2}$ and $8.5 \text{ GeV}^2 \leq Q^2 \leq 15 \text{ GeV}^2$ by the ZEUS and H1 Collaborations at HERA [1]. For larger values of Q, of relevance to the calculation of neutrino-nucleon cross section, the HERA experiments have measured $F_2^{ep}(x, Q^2)$ down to a value of $x = 2 \times 10^{-2}$. Thus, for neutrino energies above $E_\nu \sim 10^5$ GeV, the relevant small-x regime is unconstrained by experiments. In the absence of measurements in the appropriate range of (x, Q), the parton distributions are traditionally obtained by assuming certain form at some fixed value of $Q = Q_0$, i.e.

$$xq_v(x, Q_0^2) = A_c x^{\beta_v}(1-x)^{\eta_v} f_v(\sqrt{x}) \tag{2.31}$$

$$xq_s(x, Q_0^2) = A_s x^{-\lambda}(1-x)^{\eta_s} f_s(\sqrt{x}) \tag{2.32}$$

$$xg(x, Q_0^2) = A_g x^{-\lambda}(1-x)^{\eta_g} f_g(\sqrt{x}) \tag{2.33}$$

where $f(\sqrt{x})$'s are polynomials in $\sqrt{x}$ such that $f(\sqrt{x}) \to 1$ for $x \to 0$ [7]. The value for Q_0 is taken to be of the order of GeV, and the distributions are evolved to larger values of Q via QCD evolution equations. The parameters that appear in Eqs. (2.31-2.33) are determined by fitting a variety of experimental data on structure functions for $x > 10^{-3}$ and moderate values of Q [7], as well as the recent HERA data [1].

The theoretical uncertainties in the total cross section for UHE neutrinos in the standard model are therefore due to the parameterization at Q_0, the extrapolation to values of x below current data and in the evolution of the parton distribution functions to $Q > Q_0$. The issues of the parameterization of the distribution functions and the evolution are related, as we describe below.

Currently there are two theoretical approaches to understanding small-x evolution with Q, both based on perturbative QCD. The first, and more traditional approach is to determine parton densities for $Q > Q_0$ by solving the next-to-leading order Gribov-Lipatov-Altarelli-Parisi (GLAP) evolution equations numerically [8]. This is the approach used by the CTEQ Collaboration [9] and by Martin, Roberts and Stirling [10] in their parton model parameterizations which are fitted to the data in deep inelastic scattering and collider experiments. Recent HERA data at $Q^2 \sim 15 GeV^2$ [1] indicate that λ parameter that appears in Eqs. (2.31-2.33) is somewhere between 0.28 and 0.35. These values of λ are used as input in both CTEQ and MRSA distributions. Pre-HERA distribution functions, which assume $\lambda = 0$ (i.e. flat initial distributions), such as EHLQ [13] and MRS D0 [7], have been found to have large leading and next-to- leading order perturbative resummation effects [15] indicating that the fixed-order evolution with this particular assumption is a poor approximation. On the other hand, "singular" parton distributions, such as MRSA, MRS D– and CTEQ have much smaller corrections [15].

In the small-x, $x \leq 10^{-2}$ and $x >> 10^{-8}$ it is possible to approximately solve the leading order GLAP equations. The sea quark distribution is

$$xq_s(x, Q^2) = C(Q^2)\sqrt{\frac{2(\xi - \xi_0)}{\rho}} \exp\{[(2\rho(\xi - \xi_0)]^{1/2}\}, \tag{2.34}$$

and is called the double-logarithmic-approximation (DLA). Here, $\rho = (8N/b_0)\ln(1/x)$, $\xi(Q) = \ln\ln(Q^2/\Lambda^2)$, $N = 3$ is the number of colors, and $b_0 = (11N - 2n_f)/3$ for n_f flavors. This was the form of the extrapolation below $x_{\min} = 10^{-4}$ for the Eichten *et al.*[13] parton distributions that was used in the Ref. [6] calculation of total neutrino-nucleon cross section.

Recently, detailed study of the applicability of the DLA has been done by Ellis, Kunzst and Levin [14]. They showed that when λ is smaller than 0.5 for $Q \sim M_W$, the small-x solution to the GLAP equations has an x dependence

$$xq_s(x, Q^2) \sim A(Q)x^{-\lambda} \tag{2.35}$$

given a starting distribution of the same form, with $\lambda = 0.5$.

The second approach to small-x evolution is to solve the Balitskii-Fadin-Kuraev-Lipatov (BFKL) equation. This is effectively a leading $\alpha_s \ln(1/x)$ resummation of soft gluon emissions [11]. The main difference between the GLAP and BFKL approach is that the small-x behavior is the *input* (at some scale Q_0) in the GLAP approach, while the small-x behavior in BFKL approach is generated dynamically by solving the evolution equation and resumming the $\ln(1/x)$ terms. The solution to the resummed leading $\alpha_s \ln(1/x)$ evolution equation yields $\lambda \simeq 0.5$. In addition to predicting the singular behavior in x, the BFKL approach also predicts a steeper Q^2 behavior than GLAP, namely

$$xq_s(x,Q^2) \sim (Q^2)^{\frac{1}{2}} x^{-0.5}. \qquad (2.36)$$

If one assumes the singular input distributions, with $\lambda \simeq 0.5$, the Q^2 dependence obtained via GLAP evolution is $\ln Q^2$, which is slower than the BFKL prediction.

One should keep in mind, however, that the BFKL evolution breaks down for large-Q^2. The GLAP approach is applicable in the "not-so-small "-x and "large"-Q^2 region, while the BFKL solution is reliable in the small-x and moderate-Q^2 region. In case of ultra-high energy neutrino-nucleon interactions, the region of interest is small-x *and* large-Q^2, which requires a resummation of *both* $\ln 1/x$ and $\ln Q^2/Q_0^2$ contributions. Even though some progress has been made in developing a "unified" evolution equation [12], the full solution and fits to global data are far from being achieved.

Currently, however, there are several parton distribution sets obtained perturbatively using next-to-leading order GLAP evolution equations. These are applicable for the calculation of the total neutrino cross section for $E_\nu < 10^5$ GeV, and are a reasonable starting point for calculations for higher energy neutrinos. The most recent parameterizations provided by the CTEQ Collaboration and Martin, Roberts and Stirling include fits to data on lepton asymmetries in W production. In the limit of small-x, their parameterizations given by Eqs. (2.31-2.33) reduce to

$$xg(x,Q_0^2) \sim A_g x^{-\lambda} \qquad (2.37)$$

$$x(\bar{u}(x,Q_0^2)+\bar{d}(x,Q_0^2))/2 \sim A_s x^{-\lambda} \qquad (2.38)$$

The CTEQ distributions (CTEQ3) assume that $Q_0 = 1.6$ GeV and $\lambda = 0.286$, while Martin, Roberts and Stirling distributions (MRSA) have $Q_0 = 2$ GeV, and $\lambda = 0.3$. They both evolve distributions to larger Q using the next-to-leading order GLAP evolution in the modified minimal subtraction scheme ($\overline{\text{MS}}$) scheme. The λ values used by these two parametrization are thus intermediate between the two extremes that yield the small-x forms in Eqs. (2.37) and (2.38).

A more singular form of the parton distributions at small x, motivated by BFKL dynamics, appears in the MRS D− set [7]. The MRS D− parton distributions are parameterized so that in the limit of very small-x for $Q = Q_0$ the behavior is $xq_s(x,Q) = C(Q)x^{0.5}$. The power $\lambda = 0.5$ does not changes significantly with x and Q for small-x and $Q \sim M_W$. An evaluation of the numerically evolved D- distribution functions gives $\lambda \simeq 0.48$ at $x = 10^{-4}$ for Q between 50 GeV and 150 GeV. This value of λ is little changed as one goes to $x = 10^{-5}$, the limit of the numerically evolved D− parton distribution functions.

In our calculation of the total neutrino-nucleon cross section, we use the CTEQ3 parton distributions, with NLO GLAP evolution, as our canonical set. We choose the deep-inelastic scattering factorization scheme (DIS) parameterization of the parton distribution functions, which has $\lambda = 0.332$. These are labeled as CTEQ-DIS in the discussion below. The CTEQ3 distributions are particularly useful because the numerical evolution is provided for $x \to 0$, albeit, in a region where one does not expect GLAP evolution to be reliable.

To estimate the uncertainty in the total neutrino cross section, we use two extreme cases for the small-x parameterizations. To get an upper limit, we use the MRS D– set. Below $x_{\rm min} = 10^{-5}$, we extrapolate the sea quark distributions with the form of Eq. (2.38), with $\lambda = 0.48$. The normalization constant $A(Q)$ is evaluated as a function of Q at $x_{\rm min}$. The D– distributions appear to slightly overestimate the low-Q^2 HERA data [1] for $x \sim 10^{-4} - 10^{-2}$. Thus, for large E_ν, they can provide a reasonable upper limit on the cross section. To estimate a lower limit on the cross section, we use the DLA form of Eq. (2.39) for $x < x_{\rm min} = 10^{-2}$ as an extrapolation of the leading order parameterization of the CTEQ3 parton distribution functions (labeled CTEQ-DLA). To illustrate the range of parton distributions that these choices represent, we plot $x(\bar{u} + \bar{d})$ versus x for $Q = M_W$ in Ref. 2.

We calculate charged current (CC) cross sections for the three choices of parton distribution functions. The differences in the cross section appear at neutrino energies $E_\nu \sim 10^6$ GeV, reflecting the different behaviors of $xq_s(x,Q)$ below $x \sim M_W^2/(2M \cdot 10^6$ GeV$)\sim 3 \cdot 10^{-3}$. The D– and CTEQ-DIS cross sections are essentially equal for $E_\nu = 10^6$ GeV, but the CTEQ-DLA cross section is $\sim 15\%$ lower. By $E_\nu = 10^{10}$ GeV, there is a factor of 3 difference between the DLA extrapolation and the D– calculations of the cross section. The CTEQ-DIS result lies a factor of two above the DLA cross section. At $E_\nu = 10^{12}$ GeV, the discrepancy between the DLA and D– calculations is close to a factor of 10. One should keep in mind, however, that the DLA is an approximation valid for $E_\nu << 10^{11} GeV$.

We have compared our results with the recent measurement of charged-current cross section at HERA [17]. We find that our cross sections are in good agreement with the data. At $E_\nu = 100 TeV$ we find that the valence quark contribution is 15% while at energies above $10^5 TeV$ the main contribution to the total cross section comes from the sea quarks. Our results for the charged-current and neutral-current cross sections at energies $100 TeV < E_\nu < 10^4 GeV$ are 17% to 33% larger than previous Reno-Quigg results [6].

In order to calculate the number of upward-moving muons that can be detected with future neutrino detectors, such as DUMAND II [19], AMANDA [21] and NESTOR [20], we need to calculate the probability that neutrino passing on a detector trajectory creates a muon in the rock which traverses the detector,

$$P(E_\nu, E_\mu^{th}) = N_A \sigma_{CC}(E_\nu) R(E_\nu, E_\mu^{th})$$

and the attenuation of the neutrino flux in the Earth, given by

$$S(E_\nu, \Theta) = exp[-N_A(\sigma_{CC}(E_\nu) + \sigma_{NC}(E_\nu/1-y)) \int_0^{R(\Theta)} \rho(z)dz]$$

Here $R(E_\nu, E_\mu^{th})$ is the muon range and Θ is the zenith angle. We find that both the muon probability and the neutrino attenuation are *very* sensitive to the small-x behavior of the sea quark distributions, but their product is not. Thus, for the downward-moving muons (and for the contained events) the event rates obtained with MRS D- or CTEQ distributions are much larger than the same with EHLQ structure functions, but the upward rates differ only by 17% for the whole energy range.

We find that future detectors, such as DUMAND II, AMANDA and NESTOR will be able to test different models for neutrino production in the AGNs. For $E_\mu^{th} = 1TeV$, we find that DUMAND II (AMANDA, NESTOR) would be able to detect 28 (140) upward-moving muons per year originating from the diffuse AGN neutrinos, with the atmospheric background of 26 (130) events. For $E_\mu^{th} = 10TeV$, signal to background ratio becomes even better, with signal being 48 events per year on the top of 10 background events (for AMANDA and NESTOR). We plan to look into different models for production of gamma rays and neutrinos in AGN's and in the astrophysical point sources. It seems plausible that recently observed TeV gamma rays from Markarian 421 and Markarian 501 originate from both the electromagnetic interactions (inverse Compton scatterings) and the hadronic interactions. Determining the dominant mechanism for the particular energy range will be one of the problems that we will consider as well as the absorption of gamma ray in the source and by the low-energy intergalactic photons. We would like to use EGRET observations of MeV gamma rays to determine at which energies the photoproduction process in the core or in the jets of AGNs becomes significant source of the observed gamma rays. Clarifying these issues would result in better determination of the neutrino flux which is the largest uncertainty in the prediction of the event rates.

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Chapter 3

Collaborations and Outreach Activities

In addition to Professors Carruthers and Sarcevic, the grant currently supports a graduate student (Peter Valerio) and a postdoctoral fellow (Dr. Zheng Huang). Dr. Huang has joined our group since August 1995. He was previously a postdoc fellow at LBL, working with X.-N. Wang and M. Suzuki. Dr. Huang's recent work has been on the disoriented chiral condensate, η production and dilepton production in heavy-ion collisions. We anticipate a very productive collaboration with him on many topics in heavy-ion physics. Our new graduate student, Igor Zubarev, has joined our research group since this Fall. Mr. Zubarev came from the Institute for Nuclear Physics in St. Petersburg with highest recommendations from Mark Strikman (Penn State). He is currently supported by the University of Arizona. With these new additions, we shall have a stronger nuclear physics research group.

Through various arrangements we have also managed to have several visitors with partial university and partial grant support. Profs. Satz, Cleymans, W. Greiner and M. Greiner spent several weeks in Tucson in 1995. Dr. Elze has been spending this year at the University of Arizona without a support from this grant. Former students (H. Eggers and T. Hakioglu) and a former postdoc (P. Lipa and M. Greiner) continue to work closely with the Arizona core group and to extend the ideas mentioned above. The support for these activities has been quite marginal. They could not be maintained in future without additional support.

Peter Carruthers spends couple of months each year visiting Max-Planck Institute in Munich and University of Franfurt as Humboldt Fellow. Former postdoc Peter Lipa is now at the High Energy Physics Institute in Vienna. Former student Hans Eggers is now a postdoc at the High Energy Physics Institute in Vienna after one year at McGill University and one year in Regensburg as a von Humboldt Fellow. Former student Tugrul Hakioglu is a tenure track Assistant Professor at Bilkent University in Ankara, Turkey. Each continues to visit Tucson and to collaborate on a number of topics. Examples of such papers can be found bellow. We also plan to continue on-going collaborations, in particular with C. Quigg (FNAL), X.-N. Wang (LBL), M. Suzuki (LBL), R. Gandhi (Mehta Research Inst.), M.H. Reno (Iowa), H. Satz (CERN), H.-T. Elze (Regensburg) and G. Baym (Illinois).

This summer Sarcevic is organizing Workshop on "Interface of Soft and Hard Processes in QCD" at the Aspen Center of Physics, (June 3-23, 1996). This topic is of special interest to the Heavy-Ion community. The intent is to bring together experts on the theory of "hard" (high-p_T) hadronic processes in QCD (e.g. QCD jets, small-x physics, particle correlations in jets, etc.) and others who have concentrated on the description of "soft (low-p_T) hadronic processes which cannot yet be understood from first principles, but dominate hadronic and heavy-ion interactions. Theoretical work in the two areas has started to overlap, for example, with the development of a quantitative two-gluon-exchange model for the "Pomeron" in QCD by Lipatov, Mueller, and others, and conversely, with the development of alternative models in which, for example, the rise in hadronic cross sections at very high energies is driven by semihard scattering at the parton level. The development of a clear picture of how hard and soft processes merge has also become a key problem in the application of QCD to real hadronic processes. For example, the quantitative predictions for the "rapidity gaps" in hadronic interactions, now under experimental investigation and thought to be crucial for future studies of rare processes, depend on a detailed understanding of both hard and soft effects.

List of opics to be covered include:

- Lipatov Pomeron, minijets, and soft physics;
- The Pomeron and small-x problems in QCD;
- Rapidity gaps and other fluctuation phenomena;
- Parton distributions at small-x, Nuclear Shadowing Effect
- Applications to accelerator, cosmic ray, and heavy-ion physics

Partial list of potential participants:

J.D. Bjorken (SLAC), A. Mueller (Columbia), B. Muller (Duke), E. Shuryak (Stony Brook), H. Satz (CERN), M. Strikman (Penn State), J. Qiu (Iowa State), X.-N. Wang (LBL), L. Xiong (Stony Brook), J. Stachel (Stony Brook), Z. Huang (U of Arizona), Yu.L.

Dokshitzer (Lund), H.J. Lu (U of Arizona), S.D. Ellis (U.of Washington), L.L. Frankfurt (Tel Aviv) , V.N. Gribov (Budapest), F. Halzen (Wisconsin), V.A. Khoze (Durham), J. Kwiecinski (Orsay), R. Landshoff (Cambridge), E.M. Levin (Fermilab), L.N. Lipatov (St. Petersburg), N.N. Nikolaev (Julich), M.G. Ryskin (Rutherford), D. Soper (Oregon), G. Sterman (Stony Brook), C.I. Tan (Brown), C. Taylor (Case Western), W.K. Tung (Michigan State), J.Tran Thanh Van (Orsay), S. Webber (Cambridge), A. White (ANL), D. Zeppenfeld (Wisconsin).

Chapter 4

Publications and Curriculum Vitae

4.1 Ina Sarcevic

Publications 1993-1995

1. **"Jets and Jet Multiplicities in High Energy Photon-Nucleon Interactions,"** L. Durand, K. Honjo, R. Gandhi, H. Pi and I. Sarcevic, *Phys. Rev.* **D47**, 4815 (1993).

2. **"High Energy Photon-Nucleon and Photon-Nucleus Cross Sections,"** L. Durand, K. Honjo, H. Pi, R. Gandhi and I. Sarcevic, *Phys. Rev.* **D48**, 1048 (1993).

3. **"Intermittency and Multidimensional Correlations,"** I. Sarcevic, in Proceedings of the *VIIth International Symposium on Very High Energy Cosmic Ray Interactions*, L. Jones, ed. (American Institute of Physics, New York, 1993), p. 364.

4. **"Intermittency and Particle Correlations: Theory,"** I. Sarcevic, in Proceedings of the *XXVI International Conference on High Energy Physics*, J. R. Sanford, ed. (American Institute of Physics, New York, 1993), p. 919.

5. **"The Statistical Field Theory of Multiparticle Density Fluctuations,"** H. Eggers, T. Elze and I. Sarcevic, in Proceedings of the *XXII International Symposium on Multiparticle Dynamics*, C. Pajares, ed. (World Scientific, Singapore, 1993), p. 346.

6. **"Photoproduction Cross Sections at HERA Energies,"** I. Sarcevic, in Proceedings of the *Workshop on Physics at Current Accelerators and the Supercollider*, J. Hewett and A. White, eds. (World Scientific, Singapore, 1993), p. 609.

7. **"The Total and Jet Photoproduction Cross Sections at HERA and Fermilab Energies,"**, I. Sarcevic, in Proceedings of the *International Conference on Elastic and Diffractive Scatterings, (Blois'93)* (World Scientific, Singapore, 1994) p. 77.

8. **"QCD Cascades,"**, I. Sarcevic, in Proceedings of the Workshop on *Pre-Equilibrium Parton Dynamics in Heavy-Ion Collisions*, ed. X. N. Wang, LBL Report 1993.

9. **"Effective Field Theory of Multiparticle Correlations,"**, in Proceedings of the *XXIII International Symposium on Multiparticle Dynamics*, M. Block and A. White, eds. (World Scientific, Singapore, 1994) p. 287.

10. **"Three-dimensional Statistical Field Theory for Density Fluctuations in Heavy- Ion Collisions"**, H. C. Eggers, H.-Th. Elze and I. Sarcevic, *Int. J. of Mod. Phys.* **A9**, 3821 (1994).

11. **"Total and Jet Photoproduction Cross Sections at HERA,"** I. Sarcevic, in Proceedings of the *XXVIIIth Recontres de Moriond, QCD and High Energy Hadronic Interactions*, ed. J. Tran Thanh Van (Editions Frontiers, 1994), p. 535.

12. **"The Total and Jet Photoproduction Cross Sections at HERA and Fermilab Energies"**, I. Sarcevic, in Proceedings of the *Vth International Conference on Elastic and Diffractive Scattering* eds. K. M. Fried, K. Kang and C-I Tan (World Scientific, Singapore, 1994) p. 77.

13. **"Charm Production at RHIC to $O(\alpha_s^3)$"**, I. Sarcevic and P. Valerio, *Phys. Lett.* **B338** 426 (1994).

14. **"Toward the Field Theoretical Description of Multiparticle Production in High-Energy Collisions"**, in Proceedings of the *NATO Workshop on Hot Hadronic Matter*, eds. H. H. Gutbord, J. Leyessier and J. Rafelski (Plenum Pub. Co., New York, 1995) p. 283.

15. **"Charm Production at RHIC and LHC to $O(\alpha_s^3)$"**, in Proceedings of the *1994 Meeting of the American Physical Society, Division of Particles and Fields*, (DPF '94), ed. S. Seidel (World Scientific, Singapore, 1995) p. 1243.

16. **"Charm Production at RHIC and LHC to $O(\alpha_s^3)$"**, *Nucl. Phys.* **A590** (1995) 499c.

17. **"Charm Production in Hadronic and Heavy-Ion Collisions at RHIC and LHC Energies to $O(\alpha_s^3)$"**, I. Sarcevic and P. Valerio, *Phys. Rev.* **C51** 1433 (1995).

18. **"Ultrahigh Energy Neutrino Interactions"**, R. Gandhi, C. Quigg, M. H. Reno and I. Sarcevic, AZPH-TH/95-15, submitted to Astroparticle Journal.

19. **"Small-x Parton Densities from HERA and the Ultrahigh Energy Neutrino-Nucleon Cross Sections"**, R. Gandhi, C. Quigg, M. H. Reno and I. Sarcevic, AZPH-TH/95-25, presented at the *VIth International Conference on Elastic and Diffractive Scattering*, Chateau de Blois, June 20-24, 1995, to appear in *Nucl. Phys.* **B**.

20. **"Ultrahigh Energy Neutrino-Nucleon Cross Sections"**, R. Gandhi, C. Quigg, M.H. Reno and I. Sarcevic, AZPH-TH/95-38, presented at the *Fourth International Workshop on Theoretical and Phenomenological Aspects of Underground Physics, (TAUP95),*, Toledo, Spain, September 17-21, 1995, to appear in the Proceedings.

21. **"Domain Structure of a Disoriented Chiral Condensate from a Wavelet Perspective"**, Z. Huang, I. Sarcevic, R. Thews and X.-N. Wang, AZPH-TH/95-24, submitted to *Phys. Rev.* **D**.

Invited Talks 1993-1995

1. "Total and Jet Photoproduction Cross Sections at HERA", presented at *XXVIIIth Recontres de Moriond, QCD and High Energy Hadronic Interactions*, Les Arcs, France, March 20-27, 1993.

2. "Photoproduction at HERA", seminar, CERN, Theory Group, March 1993.

3. "Photoproduction Cross Sections at HERA Energies", presented at the *Workshop on Physics at Current Accelerators and the Supercollider*, Argonne National Laboratory,

Argonne, IL, June 2-5, 1993.

4. "The Total and Jet Photoproduction Cross Sections at HERA and Fermilab Energies", presented at the *International Conference on Elastic and Diffractive Scattering ("Blois '93 ")*, Brown University, Providence, RI, June 8-12, 1993.

5. "Effective Field Theory of Multiparticle Correlations", presented at the *Workshop on QCD at Finite Temperature*, Institute for Theoretical Physics, University of California, Santa Barbara, CA, August 16-23, 1993.

6. "QCD Cascades", presented at the *Workshop on Pre-Equilibrium Parton Dynamics in Heavy- Ion Collisions*, Lawrence Berkeley Laboratory, University of California, Berkeley, August 23 - September 3, 1993.

7. "Effective Field Theory of Multiparticle Correlations", presented at the *XXIII International Symposium on Multiparticle Dynamics*, Aspen, CO, September 12-17, 1993.

8. "Problems, Puzzles and Prospects: Personal Perspective on Present Particle Physics", Colloquium, Physics Department, University of Arizona.

9. "Fluctuations and Coherence in High-Energy Heavy-Ion Collisions," seminar, Nevis Labs, Columbia University, April 1994.

10. "Intermittency in High-Energy Collisions", seminar, Physics Department, State University of New York at Stony Brook, April 1994.

11. "Intermittency in High-Energy Collisions: Recent Developments", seminar, Physics Department, Brookhaven Laboratory, April 1994.

12. "Toward the Field Theoretical Description of Multiparticle Production in High-Energy Collisions", presented at the *NATO Workshop on Hot Hadronic Matter*, Divonne-les Bains, France, June 27- July 1, 1994.

13. "Charm Production at RHIC and LHC to $O(\alpha_s^3)$", presented at the *1994 Meeting of the American Physical Society, Division of Particles and Fields, (DPF'94)*, Albuquerque, August 2-6, 1994.

14. "Charm Production at RHIC and LHC to $O(\alpha_s^3)$", presented at *Quark Matter '95*, Monterey, California, January 9-13, 1995.

15. "The Ultrahigh Energy Neutrino Interactions", seminar, Physics Department, UCLA, Los Angelos, May 1995.

16. "Ultrahigh Energy Neutrino Interactions: How to Search for AGN Neutrinos", seminar, Theoretical Division, T-8, Los Alamos National Laboratory, Los Alamos, June 1995.

17. "Open Charm Production in Hadronic and Heavy Ion Collisions at RHIC and LHC Energies to $O(\alpha_s^3)$", Nuclear and Particles Colloquium, Los Alamos National Laboratory, Los Alamos, June 1995.

18. "Ultrahigh Energy Neutrino Interactions: Implications for AGN Neutrino Searches", seminar, Astronomy Department, UCLA, Los Angelos, June 1995.

19. "Ultrahigh Energy Neutrino Interactions", seminar, LPTHE Universite Paris XI, Orsay, France, June 1995.

20. "Small-x Parton Densities from HERA and the Ultrahigh Energy Neutrino-Nucleon Cross Sections", presented at the *VIth International Workshop on Elastic and Diffractive Scatterings, Frontiers in Strong Interactions*, Recontres de Blois, Chateaue de Blois, France, June 1995.

21. "Small-x Parton Densities from HERA and the Ultrahigh Energy Neutrino-Nucleon Cross Sections", presented at the *Fourth International Workshop on Theoretical and Phenomenological Aspects of Underground Physics, (TAUP95)*, Toledo, Spain, September 1995.

22. "Particle Correlations and the Disorentied Chiral Condensate from a Wavelet Perspective", invited talk to be presented at *XXIV Coral Gables Conference on Elementary Particle Physics and Cosmology*, Coral Gables, Florida, January 1996.

4.2 Peter Carruthers

Selected Activities

1. Guest professor, Institute for Theoretical Physics, University of Frankfurt, March 7–July 7 1995. The salary and air fare were paid from a Humboldt Prize and the Gesellschaft für Schwerionenforschung (GSI) in Darmstadt.

2. Honorary Trustee, Aspen Center for Physics.

3. Universities Research Association, President's Council, proxy for University of Arizona President M. Pacheco.

4. Associate Editor of the new journal "Complexity".

Recent Publications

1. H.C. Eggers, P. Lipa, P. Carruthers and B. Buschbeck,
"Higher order pion interferometry. Moments and cumulants",
Phys. Lett. **301B**, 298 (1993).

2. H.C. Eggers, P. Lipa, P. Carruthers and B. Buschbeck,
"Integral correlation measures for multiparticle physics",
Phys. Rev. D **48**, 2040 (1993).

3. M. Greiner, P. Lipa and P. Carruthers,
"Wavelet correlations in the p-model",
Phys. Rev. **E51**, 1948 (1995).

4. M. Greiner, J. Giesemann, P. Lipa and P. Carruthers,
"Wavelet correlations in hierarchical branching processes",
University of Arizona preprint AZPH-TH/94-25 (submitted to Z. Phys. C).

5. P. Lipa, M. Greiner and P. Carruthers,
"Wavelet Analysis of Multiparticle Correlations",
in Proceedings of Cracow Workshop "Soft Physics and Fluctuations", eds. A. Białas, K. Fiałkowski, K. Zalewsky, and H. C. Hwa (World Scientific, Singapore, 1994), p.105

6. P. Lipa, H. C. Eggers, P. Carruthers and B. Buschbeck,
"Measurement of Correlations in Multiparticle Production",
in Proceedings of XXIII International Symposium on Multiparticle Dynamics, Eds. M.M. Block and A.R. White, World Scientific (1994),p.268

7. P. Carruthers, M. Greiner and P. Lipa,
"Ideas on Wavelet Analysis Applied to High Energy Physics",
in Proceedings of XXIII International Symposium on Multiparticle Dynamics, Eds. M.M. Block and A.R. White, World Scientific (1994), p.281

8. H.-T. Elze and P. Carruthers,
"About Entropy and Thermalization"
in: NATO Advanced Study Workshop on Hot Hadronic Matter: Theory and Experiment, 1994; Eds. Jean Letessier, Hans H. Gutbrod and Johann Rafelski, NATO ASI Series, Plenum Press, New York (1995), p.98

9. P. Lipa, M. Greiner and P. Carruthers,
"Analysis of Multiparticle Correlations and the Wavelet Transform",
in: NATO Advanced Study Workshop on Hot Hadronic Matter: Theory and Experiment, 1994; Eds. Jean Letessier, Hans H. Gutbrod and Johann Rafelski, NATO ASI Series, Plenum Press, New York (1995), p.132

10. P. Carruthers,
Keynote speech on "Complexity, Correlations and Fluctuations",
in: Eighth International Congress on Recent Progress in Many-body Theories, Schloss Seggau, Styria, Austria, August 22-26, 1994, University of Arizona preprint AZPH/94-33 (Plenum, to be published).

11. M. Greiner, P. Lipa and P. Carruthers,
"Wavelets: A new tool to study selfsimilarity and clustering in complex reactions",
in: CORINNE II: International Workshop on Multiparticle Correlations and Nuclear Reactions, Nantes, France, Sept. 1994; (to be published in the conference proceedings).

12. P. Carruthers,
"Some Thoughts on the Irregular Galaxy Distribution in the Universe",
Saha Symposium, Calcutta, Dec. 1993, (World Scientific, to be published)

13. P. Carruthers,

"Some Thoughts on the Negative Binomial Count Distribution and Correlations", submitted to J Mod Phys A

14. P. Carruthers,
"A Personal Recollection of Count Statistics and Especially the Negative Binomial Distribution", to be published

15. P. Lipa, M. Greiner and P. Carruthers,
"Introduction to wavelets, correlations in simple branching processes and applications of wavelets to correlation measurements",
Invited Talk at the School and Workshop on Heavy Ion Collisions, Bratislava, Slovakia, September 4-9,1995. (To be published in Acta Phys. Slovakia)

16. M. Greiner, P. Lipa, J. Giesemann and P. Carruthers,
"Wavelet correlations in hierarchical branching models",
Joint writeup of three talks presented by M. Greiner, P. Lipa and P. Carruthers at the XXV International Symposium on Multiparticle Dynamics, Stará Lesná, Slovakia, 12–17 September 1995.

Recent Meetings Attended

1) "Nucleus-Nucleus Collisions V", Taormina, Italy 30 May-4 June

2) Paris "Quantum Infrared Scattering at High Energies" June 6-19.

3) Aspen Center for Physics, August 9–14

4) Keystone, Colorado,"Scientist to Scientist Colloquium": August 9-19

5) Alexander von Humboldt-Stiftung: Symposium of the German- American Academic Council, Madison, Sept. 29-30

6) VIII International Conference on Recent Progress in Many-Body Theories, Schloss Seggau, Styria, Austria, August 22-27

International Advisory Committees

1) XXIV International Symposium on Multiparticle Dynamics. Vietri sul Matri, Italy (Sept.13-18, 1994)

2) Uzgorod (Ukraine)Sept. 7-11

3) Moscow (Zvenigorod) Sept 16-22 IX International Workshop on High Energy Field Theory

Collaborators:

1. Peter Lipa (Tucson and Vienna: held postdoctoral position in Arizona for 3 and 1/2 years),
2. Hans Eggers (presently on a fellowship in Vienna; permanent position in South Africa from January 1996; got his PhD with me),
3. Ludwig Neise (now at Heidelburg University; was Lynen fellow and Adjunct Assistant Professor in Physics here for two years.)
4. Martin Greiner (Giessen University, moving to Dresden), was Lynen Fellow and Adjunct Assistant Professor in Physics here for two years.
5. H.–Thomas Elze (CERN). He has a Heisenberg Fellowship and spends several months in Tucson each year.

4.3 Zheng Huang

Experiences

Aug, 1995 – Present	Postdoc Fellow at University of Arizona
Sept, 1993 – Aug, 1995	Postdoc Fellow at Lawrence Berkeley Lab
May, 1989 – Sept, 1993	Ph.D. Graduate at Simon Fraser University
Sept, 1985 – May, 1989	M.Sc. Graduate at Inst. of High Energy Physics
Sept, 1981 – Sept, 1985	B.Sc. Undergraduate at Pekin University

Publications 1992-1995

1. Z. Huang, I. Sarcevic, R. Thews and X.-N. Wang, *Domain Structure of Disoriented Chiral Condensate from a Wavelet Perspective*, AZPH-TH/95-24, LBL-38001 (1995), submitted to Phys.Rev.**D**.

2. Z. Huang and X.-N. Wang, *Partial $U(1)_A$ Restoration and η Enhancement in High Energy Collisions*, LBL-37539 (1995), Phys. Rev.**D**, in press.

3. Z. Huang and M. Suzuki, *General Classical Solutions of Nonlinear σ Model and Pion Charge Distribution of Disoriented Chiral Condensate*, LBL-37524 (1995), Phys. Rev. **D**, in press.

4. Z. Huang, *Thermal Dilepton Production in Heavy-Ion Collisions*, Phys. Lett. **B361**, 131 (1995).

5. Z. Huang and M. Suzuki, *Isospin Coherence and Final-State Quantum Scattering of Disoriented Chiral Condensate*, Phys. Rev. **D52**, 2610 (1995).

6. Z. Huang, *Do Longitudinal W Bosons and Higgs Boson form a Bound State?*, LBL-36910 (1994), in Proceedings of 4th International Conference on Physics Beyond the Standard Model, eds. J. Gunion, T. Han and J. Ohnemus (World Scientific, Singapore, 1995), p.585.

7. T. Han, Z. Huang and P.Q. Hung, *A Model for the Low-lying States in Strongly Interacting Electroweak Sector*, LBL-36270 (1995), Submitted to Phys. Rev. Lett.

8. Z. Huang, *Disoriented Chiral Condensate*, LBL-36740 (1994), in Proceedings of 4th International Conference on Physics Beyond the Standard Model, eds. J. Gunion, T. Han and J. Ohnemus (World Scientific, Singapore, 1995), p.422.

9. M. Asakawa, Z. Huang and X.-N. Wang, *How Disoriented Chiral Condensates Form: Quenching vs. Annealing*, Phys. Rev. Lett. **74** (1995) 3126-3129.

10. Z. Huang, M. Suzuki and X.-N. Wang, *Signals of Disoriented Chiral Condensate*, Phys. Rev. **D50** (1994) 2277

11. Z. Huang and X.-N. Wang, *Cluster Structure of Disoriented Chiral Condensate in Rapidity Distribution*, Phys. Rev. **D49** (1994) 4335-4338

12. Z. Huang, *Vacuum Misalignment in High Energy Collisions*, Phys. Rev. **D49** (1994) 16-19

13. Z. Huang, *CP Violation, Anomalous Discrete Symmetry and Dynamical Fermion Mass*, Ph.D. Thesis (1993)

14. Z. Huang and K.S. Viswanathan, *Self-mass for Massive Quark*, Phys. Rev. **D48** (1993) 4392-4397

15. Z. Huang, *Anomalous Discrete Symmetry and Spontaneous CP Violation*, Phys. Rev. **D46** (1992) 4818-4821

16. Z. Huang, *The Measure of the Strong CP Violation*, Phys. Rev. **D48** (1993) 270-282

17. Z. Huang, K.S. Viswanathan and D.D. Wu, *Comments on the Vacuum Orientations in QCD*, Mod. Phys. Lett. **A7** (1992) 3147-3154.