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NOVEL QCD EFFECTS IN NUCLEAR COLLISIONS*

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Abstract

Heavy ion collisions can provide a novel environment for testing fundamental dynamical processes in QCD, including minijet formation and interactions, formation zone phenomena, color filtering, coherent co-mover interactions, and new higher twist mechanisms which could account for the observed excess production and anomalous nuclear target dependence of heavy flavor production. The possibility of using light-cone thermodynamics and a corresponding covariant temperature to describe the QCD phases of the nuclear fragmentation region is also briefly discussed.

1. Introduction

One of the most important goals of relativistic nuclear collisions is to probe the extraordinarily rich and diverse phenomenology of quantum chromodynamics. The nucleus, in fact, provides one of the few ways we can change the physical environment in QCD. In the case of atomic physics, one can use external Stark and Zeeman electromagnetic fields to modify atomic wavefunctions and to probe the underlying dynamics. Analogously in QCD, we can study the dependence of reactions on the parameters of a nuclear medium in order to probe hadronic substructure and dynamics.

There are a number of intriguing nuclear effects which have already been identified in deep inelastic lepton-nucleus reactions and other hard scattering processes involving nuclear targets [1]. These include:

- Non-additive contributions to nuclear structure functions, such as shadowing and anti-shadowing which reflect the coherence of multiple scattering interactions of quarks and gluons within the nucleus.
- Energy-loss mechanisms and collision broadening of transverse momentum distributions caused by the elastic and inelastic interactions of quark and gluons as they propagate through a nucleus.
- Formation zone phenomena which provide a fundamental quantum mechanical limit on multiple inelastic reactions of particles in the nucleus at high energies.
- Color transparency phenomena [2,3] in quarkonium production, hard quasi-elastic reactions, and diffractive jet production in which the nucleus acts as a differential "color filter" to separate Fock components of different transverse size in the projectile's wavefunction as well as to identify perturbative short–distance subprocesses versus non–perturbative mechanisms. Conversely, "color opacity" can be used to relate high multiplicity fluctuations in particle production distributions to the large size fluctuations of the Fock wavefunctions [4].

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- Higher twist processes arising from multi-parton coherence in the hadron wavefunction. Such processes often dominate reactions at large momentum fraction x.
- Short-range correlations within the nucleus will produce particles at large $x > 1$ beyond the range allowed from single-nucleon kinematics. The observed distributions in the target fragmentation region are projectile independent and appear to be much harder than expected from usual Fermi motion calculations. These "cumulative effects" [5] can be used to identify short-range correlations within the nuclear wavefunction such as hard internal nucleon-nucleon [5] or quark-gluon interactions [6] ("intrinsic hardness") as well as "hidden color" components of the nuclear wavefunction orthogonal to the dominant color-singlet cluster components [1].
- Co-mover interactions between produced quarks and gluons and the low transverse momentum spectator partons created in the nuclear collisions. Such interactions provide an alternative mechanism for jet hadronization compared to the usual mechanisms which occur in e^+e^- annihilation. The coalescence of the produced heavy quarks with spectators moving at the same rapidity [3,7,8,9] can in fact account for the observed suppression [10] of the ratio of J/ψ to continuum $\mu^+\mu^-$ pairs seen in central nucleus-nucleus collisions, as well as the suppression of J/ψ and Υ production seen at negative x_F in proton–uranium collisions [11].
- Heavy quark and hadron production in nuclei due to the freeing of intrinsic heavy quark components in the hadronic or nuclear Fock wavefunctions [12].
- The production of nuclear-bound quarkonium, novel bound states such as $\eta_c He^3$ due to the QCD van der Waals potential [13].
- Large momentum transfer exclusive nuclear reactions such as electron-deuteron scattering and deuteron photo-disintegration which reflect multi–quark scattering processes in the nucleus and possible hidden–color components in the nuclear wavefunction.

Each of these effects will have their role amplified in nucleus–nucleus collisions. In addition to their intrinsic interest, it is clear that we need to have a systematic understanding of $A_1 + A_2$ collisions from the lightest to the heaviest heavy nuclei in order to unravel specific effects which could be attributed to a quark-gluon phase of QCD. In this survey of nuclear QCD phenomena, I will emphasize the importance of studying the complete kinematic range of relativistic heavy ion collisions, particularly the nuclear fragmentation region. I also briefly discuss the possibility of using light-cone thermodynamics and a corresponding covariant temperature to describe QCD phases of the nuclear fragmentation region.

2. Limitations of The QCD Factorization Theorem

The cornerstone of virtually all predictions for inclusive reactions in high energy physics is the factorization theorem of perturbative QCD which separates the perturbatively–calculable hard–scattering quark and gluon dynamics from the non-perturbative bound state dynamics contained in the process–independent structure functions $G_{a/A}(x, Q^2)$ and jet fragmentation functions $D_{H/a}(z, Q^2)$. For example, to leading order in $1/\mathcal{M}^2$, the inclusive cross section to produce a heavy quark pair at invariant mass squared $\mathcal{M}^2 = x_a x_b s$ and Feynman longitudinal momentum fraction $x_f = x_a - x_b$ has the form

$$\frac{d\sigma_{A_1A_2\to Q\bar{Q}X}}{dx_a dx_b} = \sum_{ab\to Q\bar{Q}} G_{a/A_1}(x_a, \mathcal{M}^2) G_{b/A_2}(x_b, \mathcal{M}^2) \widehat{\sigma}_{ab\to Q\bar{Q}}(\mathcal{M}^2 = x_a x_b s, \alpha_s(\mathcal{M}^2)).$$

This factorization holds for all projectiles and targets A_1 and A_2, leptons, photons hadrons and nuclei. Here x_a is the boost–invariant light-cone momentum fraction of the interacting partons $x_a = k_a^+/p_{A_1}^+ = (k_a^0 + k_a^z)/(p_{A_1}^0 + p_{A_2}^z)$. In the case of the Drell-Yan process $A_1A_2 \to \mu^+\mu^- X$, factorization was proved to leading order in $1/Q^2$ by Bodwin and by Collins, Soper, and Sterman and then extended to next-to-leading twist by Qiu and Sterman [14].

One of the most striking applications of QCD factorization is its predictions for high mass lepton-pairs and high transverse momentum jet production in heavy ion collisions. Since the nuclear structure functions are approximately linear in nucleon number A, the predicted rate for hard scattering reactions is approximately proportional to A_1A_2 whereas the total inclusive cross section only increases as the total radius squared, or $(A_1^{1/3}+A_2^{1/3})^2$. Thus the number of lepton pairs per interaction in uranium-uranium collisions compared to that for $p-p$ collisions is $A_U^{4/3} \sim 1500$. If we extend this result to "minijet" production, jets with $p_T > 2$ GeV, say, then one predicts a large inclusive cross section and a large number of jets per nucleus-nucleus collisions which can dominate hadron production in the central region [15]. In addition to the usual fragmentation processes as seen in SPEAR e^+e^- data, one also needs to consider the multiple scattering, energy losses, and coalescence of the minijets with co-moving spectator partons. Clearly this phenomena has to be taken into account when considering signals for QCD phase transitions in relativistic heavy ion collisions.

It is crucial to understand the physics and range of validity of QCD factorization when applying it to nuclear processes. At first sight, it seems remarkable that the QCD factorization formula could be valid for heavy nuclear projectiles or targets, considering that each incident and final state parton will multiply–scatter both elastically and inelastically as it propagates through the nucleus. Nevertheless, QCD factorization predicts that the projectile probability distribution $G_{a/A_1}(x_a, Q^2)$ must be independent of the size of the target A_2. There are two reasons why multiple scattering does not affect the inclusive cross section to leading order in $1/Q^2$: first of all, the multiple scattering of the projectile partons in the target (or *vice versa*) can cause a finite collision broadening of the parton transverse momentum distribution which will spread out the Q_T distribution of the outgoing pair with nucleon number A. However, this effect does not change the integrated cross section $d\sigma/dQ^2$ to leading order in $1/Q^2$ [16]. Nuclear broadening of the muon pair transverse momentum distribution has been reported by the NA-10 [17] and E772 Collaborations [11].

The second essential ingredient in the proof of QCD factorization is the fact that the mean energy loss dE/dz per unit length in the target, suffered by a projectile parton due to multiple elastic or inelastic scattering, is constant and independent of the projectile energy. In fact, the mean energy loss of the parton in the target is limited by quantum mechanics [18]. The basic condition is set by the uncertainty principle, $\Delta p_z \Delta z > 1$ where Δp_z is the minimum momentum transfer of the parton in an inelastic reaction which occurs within Δz, the distance between the scattering centers. The change in longitudinal momentum of the scattered parton due to induced radiation is $\Delta\mathcal{M}^2/2E_{lab}$ where $\Delta\mathcal{M}^2$ is the difference between the incident parton mass squared and the mass squared of the parton–gluon system after radiation. Thus induced radiation effectively cannot recur in a nucleus of length L_A if the parton has energy $E_{lab} > \frac{1}{2}\Delta\mathcal{M}^2 L_A$. In particular, if a soft gluon is emitted with momentum fraction x_g and transverse momentum $k_{\perp g}$, then $\Delta\mathcal{M}^2 \sim k_{\perp g}^2/x_g$, and only gluons with $x_g < k_{\perp g}^2 L_A/2E_{lab}$ can be radiated in the nucleus from a parton with energy E_{lab}. Thus the maximum energy loss of an incident or outgoing parton a participating in a high momentum transfer reaction due to induced radiation in the nucleus is of the form

$$\Delta x_a = \kappa A^{1/3}/x_a s,$$

since $E^a_{lab} = x_a s/2M_P$ and the nuclear radius $R_A \sim (1.2 \text{ fm})A^{1/3}$ characterizes the largest effective distance between scattering centers in the nucleus. The constant κ is independent of the scale Q^2 of the hard collision. At asymptotic energies, $s \to \infty$, the fractional energy loss Δx_a vanishes and the structure function of the projectile becomes independent of the target, as required by the QCD factorization theorem.

The factorization theorem breaks down in nuclear targets when the energy loss becomes significant compared to the parton energies – *i.e.* there is a target length condition $L < 2E_{lab}/\Delta\mathcal{M}^2$ for the validity of the leading twist predictions [16], where $\Delta\mathcal{M}^2$ characterizes the change in invariant mass squared of the interacting state in the inelastic collision. The physics of the target length condition is closely related to the formation zone phenomena of Landau, Pomeranchuk, and Migdal. As shown in Ref. [16] the cancellation of induced inelastic radiation from a high energy parton participating in a hard reaction occurs because of the destructive interference of the radiation emitted by the parton at the two scattering centers. For example, consider the inelastic interactions of an antiquark due to induced gluon emission in the target before it annihilates to produce a massive lepton pair. When one analyzes such processes in time-ordered perturbation theory, one can consider different on-shell intermediate state processes corresponding to which target center the inelastic reaction occurs on. However, at high momentum transfer the minimum momentum transfer Δp_z to the nucleus is so small that these classical radiation processes lead to the same final state and interfere destructively.

Let us assume that the particles radiated by the incident or final state partons have a characteristic transverse momentum squared $< k_\perp^2 >\sim 0.1 \text{ GeV}^2$. Then the energy loss constant is $\kappa \sim (1.2 \text{ fm})M_p < k_\perp^2 >\sim 0.5 \text{ GeV}^2$. This corresponds to an energy–independent mean energy loss per unit length $dE_{inelastic}/dz \sim< k_\perp^2 > /2 \sim 0.3$ GeV/fm. A similar degradation of energy is expected from elastic scattering [15]. The total predicted energy loss, $dE/dz \sim 0.6$ GeV/fm, appears to be consistent with the magnitude determined by Gyulassy *et al.* [15] using combined SLAC and EMC data for jet fragmentation in nuclei. The energy loss of an initial or final state parton in the nucleus will create extra hadronic energy in the central rapidity region. Since this effect is independent of the projectile energy, it becomes insignificant at high energies, thus explaining the lack of nuclear target dependence of jet fragmentation processes in deep inelastic lepton scattering. The same formation zone considerations of course apply to jets produced in heavy ion collisions.

The data for J/ψ production by protons and pions in nuclei show a strong attenuation at large x_F [19]. This effect does not depend on the target momentum fraction x_2 so it cannot be due to shadowing of the nuclear target quark or gluon structure functions. Recently Gavin and Milana [20] have made the interesting suggestion that this nuclear dependence could be caused by the energy loss of the incident gluon and outgoing charm quarks as they propagate through the nucleus. However, the quantum mechanical bound for the energy loss dE/dz appears to be too small to explain this effect. In any event, the parton energy loss is independent of the pair mass $\mathcal{M}^2$; thus one cannot explain the difference between the observed A-dependence of Υ and J/ψ production seen by E772 [11] at high x_F by a parton energy loss mechanism alone. As emphasized by Badier *et al.*, [19] the fusion mechanism for charm production (assuming the usual gluon momentum distributions) leaves an unexplained excess of J/ψ production at large x_F in the same region where the cross section breaks leading twist–factorization and has a $A^{2/3}$-nuclear dependence. We shall argue that the excess quarkonium production at large x_F is due to an intrinsically higher-twist mechanism where the production reaction occurs on the nuclear surface. We return to this topic in section 5.

It should be emphasized that the target–independence of jet hadronization and the factorization of the fragmentation function $D_{H/q}(z, Q^2)$ assumes that the jet is produced at large transverse momentum in a kinematic region clear of the spectator partons or hadrons created from the fragmentation of the beam and target. In QED, Bethe-Heitler pair production in coherent nucleus-nucleus collisions is strongly modified by the attractive forces between the produced electron and the highly charged ions. The Coulomb forces are particularly strong when the charges have similar velocities, leading to binding of the electron with the co-moving nucleus [21]. Similarly in QCD, a quark can hadronize by coalescing with a co-moving spectator parton [22]. The effect is clearly enhanced in nucleus-nucleus collisions and in events which have accompanying large transverse energy E_T. As discussed in Refs. [7] and [8], co-mover interactions can account for the observed suppression [10] of the J/ψ to continuum $\mu^+\mu^-$ ratio seen in nucleus-nucleus collisions at high E_T as well as the suppression of J/ψ and Υ production seen at negative x_F in proton–uranium collisions [11]. The co-moving partons will bind to the produced heavy quark before it can bind into heavy quarkonium. Just as in photon-induced capture reactions the effect could be non-linear in the number of co-moving spectators. A comprehensive discussion of the systematics of charm production based on fusion processes, intrinsic charm, and co-mover interactions, including detailed predictions for heavy quarkonium suppression from co-mover interactions in relativistic heavy ion collisions is given in Ref. [7].

3. Shadowing and Anti–Shadowing of Nuclear Structure Functions

One of the most important manifestations of nuclear effects in QCD is the shadowing and anti–shadowing of deep inelastic nuclear structure functions; i.e, the depletion of the effective number of nucleons F_2^A/F_2^N at low $x \lesssim 0.1$, and the increase above nucleon additivity at $x \sim 0.15$. Results from the EMC collaboration [23] and SLAC [24] indicate that the effect is roughly Q^2–independent; *i.e.* shadowing is a leading twist effect in the operator product analysis. In contrast, the shadowing of the real photo–absorption cross section due to intermediate vector mesons falls away as an inverse power of Q^2.

In general, shadowing of nuclear cross sections can be attributed to destructive interference of the multiple scattering amplitudes in the nucleus. In the case of a hadron-nucleus cross section the incident hadron scatters elastically on a nucleon N_1 on the front face of the nucleus. At high energies the phase of the amplitude is imaginary. The hadron then propagates through the nucleus to nucleon N_2 where it interacts inelastically. The accumulated phase of the hadron propagator is also imaginary, so that this two–step amplitude is coherent and opposite in phase to the one–step amplitude where the beam hadron interacts directly on N_2 without initial–state interactions. Thus the target nucleon N_2 sees less incoming flux: it is shadowed by elastic interactions on the front face of the nucleus. If the hadron–nucleon cross section is large, then for large A the effective number of nucleons participating in the inelastic interactions is reduced to $\sim A^{2/3}$, the number of surface nucleons.

In the case of virtual photo–absorption, the photon converts to a $q\bar{q}$ pair at a distance before the target proportional to $\omega = x^{-1} = 2p \cdot q/Q^2$ in the laboratory frame. In a physical gauge, such as the light–cone $A^+ = 0$ gauge, the final–state interactions of the outgoing parton can be neglected in the valence approximation, and effectively only one member of the quark–antiquark pair interacts. The nuclear structure function F_2^A producing quark q can then be written as an integral [25,26] over the inelastic cross section $\sigma_{\bar{q}A}(s')$ where s' grows as $1/x$ for fixed space–like anti–quark mass. Similarly, the anti–quark nuclear structure function is related to inelastic quark–nucleus scattering. Thus the A–dependence of the deep inelastic nuclear structure functions reflects the A–dependence of the q and $\bar{q}$ cross sections in the nucleus.

Hung Jung Lu and I have recently applied the standard Glauber analysis to $\sigma_{\bar{q}A}$ and σ_{qA} assuming that formalism can be taken over to off–shell interactions [27]. In analogy to hadron–nucleon interactions, the coupling of the pomeron to the quark-nucleon scattering amplitude leads to shadowing of the quark-nucleus interactions at high energy and thus shadowing of the nuclear structure functions at low x. The magnitude of shadowing predicted by the model is consistent with the data for $0.1 > x > 0.01$; below this region, one expects higher–twist and vector–meson dominance shadowing to contribute. More surprisingly, the multi–scattering analysis also predicts anti–shadowing of the structure functions. Anti–shadowing requires the presence of kinematic regions where the real part of the $\bar{q} - N$ amplitude dominates over the imaginary part. In fact, if one introduces an $\alpha_R \simeq \frac{1}{2}$ Reggeon contribution to the $\bar{q}N$ and qN amplitudes, the real phase introduced by such a contribution leads to constructive interference in the quark-nucleus cross section at intermediate energies and thus anti–shadowing in the nuclear structure functions at $x \simeq 0.15$ at the few percent magnitude seen by the SLAC and EMC experiments [23,24]. It should be emphasized that the constructive interference which gives anti–shadowing in the $x \sim 0.15$ region is automatic since the phase of the Reggeon term is dictated by analyticity from the power behavior $F_2(x) \propto x^{1-\alpha_R}$ of the non-singlet structure functions at low x.

An advantage of the multi-scattering analysis is that it correlates shadowing phenomena to microscopic quark–nucleon parameters. The parameters for the effective quark–nucleon cross section required to understand shadowing phenomena provide important information on the interactions of quarks in nuclear matter. This approach also provides a dynamical and analytic explanation of anti–shadowing which was originally predicted [28] on the basis of conservation laws. The multi–scattering analysis also provides the input or starting point for the logarithmic evolution of the nuclear deep inelastic structure functions, as given by Mueller and Qiu [29]. Using the perturbative QCD factorization theorem for inclusive reactions, the same analysis can be extended to Drell–Yan and other fusion processes, taking into account the separate dependence on the valence and sea quarks. In the case of hard scattering reactions in nucleus-nucleus collisions QCD factorization predicts analogous shadowing and anti-shadowing phenomena in both the target and projectile.

4. The Nucleus as a QCD Filter

There are a large number of ways in which a nuclear target can probe fundamental aspects of QCD. A primary concept is that of the "color filter" [30,31]: if the interactions of an incident hadron are controlled by gluon exchange, then the nucleus will be transparent to those fluctuations of the incident hadron wavefunction which have small transverse size. Such Fock components have a small color dipole moment and thus will interact weakly in the nucleus; conversely, Fock components of normal hadronic size will interact strongly and be absorbed during their passage through the nucleus [30]. For example, large momentum transfer quasi–exclusive reactions [32], are controlled in perturbative QCD by small color–singlet valence–quark Fock components of transverse size $b_\perp \sim 1/Q$; thus initial–state and final–state corrections to these hard reactions are suppressed, and they can occur in a nucleus without initial or final state absorption or multiple scattering of the interacting hadrons. Thus, at large momentum transfer and energies, quasi–elastic exclusive reactions are predicted to occur uniformly in the nuclear volume. This remarkable phenomenon is called "color transparency" [2]. Thus QCD predicts that the cross section for quasi–elastic proton-proton scattering in nucleus-nucleus collisions will be proportional to the product of proton numbers in the beam and target:

$$\frac{\frac{d\sigma}{dt}[A_1A_2 \to pp(A_1-1)(A_2-1)]}{\frac{d\sigma}{dt}(pp \to pp)} \to Z_1Z_2$$

at large momentum transfer squared t. The signal for these events will be two nearly coplanar protons, with no other hadrons produced. The cross section is integrated over Fermi motion and the excited states of the spectator nuclei. In contrast to the QCD color transparency prediction, the traditional theory of nuclear absorption predicts that quasi–elastic scattering occurs primarily on the front surface of the nucleus. The above ratio would the be proportional to $Z_1^{2/3}Z_2^{2/3}$, *i.e.* the number of protons exposed on the nuclear surfaces. The kinematic conditions for the validity of the color transparency predictions are discussed in Ref. [33].

There are, however, indications that QCD transparency fails to hold in the case of pA collisions at $s_{NN} \sim 5\ GeV^2$ [34]. de Teramond and I have noted [35] that this anomaly could be due to charm production in the intermediate state of the pp elastic amplitude. The same effect can also account for the anomalously large spin-spin correlations seen in large angle pp scattering at the same energy.

5. Anomalous Production at Large x

One of the few areas of high energy physics where data appear to be in strong conflict with the predictions of the leading–twist QCD factorization theorem is charm hadroproduction. In particular, the data [19, 36, 11] on high momentum J/ψ production in hadron-nucleus collisions exhibit a number of remarkable results:

1. When the nuclear number (A) dependence of the J/ψ production cross section is parametrized as A^α, the effective power α decreases with the fractional momentum $x = x_F$ of the J/ψ; *i.e.* the data at different energies display Feynman scaling [37], $\alpha = \alpha(x)$. Contrary to expectations, the power α is not a function of x_2, the fraction of momentum taken from the nucleus and given to the J/ψ. This result violates the perturbative QCD factorization prediction for hard processes, and thus its explanation must lie in higher twist effects.
2. The cross section measured by NA3 [19] for J/ψ production by protons at large x appears to have a "diffractive contribution" in excess of what is predicted from conventional $gg \to c\bar{c}$ and $q\bar{q} \to c\bar{c}$ fusion subprocesses.
3. The transverse momentum distribution of the J/ψ produced in πW collisions significantly narrows at high x [38].
4. The J/ψ is normally produced with no net polarization. However, at $x > 0.9$ the J/ψ produced in pion-Tungsten collisions are almost completely longitudinally polarized [38]. This fact alone strongly suggests that a new production mechanism for heavy quarkonium must dominate in the limit $x \to 1$.

In addition to the anomalies seen in charmonium production, a number of experiments have reported anomalously flat open charm hadroproduction at large x, particularly measurements of charmed baryons [39,42]. None of these anomalies can be readily explained by conventional leading twist $gg \to c\bar{c}$ or $q\bar{q} \to c\bar{c}$ fusion subprocesses. It should also be noted that the EMC measurement [40,41] of the charm structure function of the nucleon appears to be much larger than expected from photon-gluon subprocesses at large x_{Bj}.

There is, however, another source for heavy quark production in QCD. From the perturbative point of view, a $uudc\bar{c}$ Fock component can also be generated by the $gg \to c\bar{c}$ amplitude where the gluons are emitted from two or more of the valence quarks. The probability for finding the heavy quark pair of mass $M_{Q\bar{Q}}$ or greater in the hadron wavefunction from this mechanism is thus of order $\alpha_s^2(M_{Q\bar{Q}}^2)/M_{Q\bar{Q}}^2$, with the overall coefficient set by the parton-parton correlation length. Intrinsic charm is thus a higher twist mechanism. The coupling of the charm pair to more than one valence quark in the proton also implies that more of the parent hadron's momentum is carried by the heavy quarks. From the non-perturbative point of view, the intrinsic charm Fock components are associated with the bound-state equation in which constituents tend to have equal velocity. Thus unlike normal sea quarks generated by evolution, the heaviest constituents, the intrinsic charm quarks, will tend to have a large fraction of the parent hadron's momentum [43]. If the excess in the measured charm structure function is identified with intrinsic charm, then one requires a 0.3% probability for the intrinsic charm Fock state in the nucleon [41]. Heavier quark pairs in the proton wavefunction will be suppressed inversely as the square of their mass [44].

If the projectile wavefunction has an intrinsic charm contribution [43], then one can readily account for the production of heavy quarks at large x_F as well as the observed Feynman scaling of the nuclear dependence. For example, in a hadronic collision the intrinsic c and $\bar{c}$ can coalesce to produce a charmonium state with the majority of the projectile's momentum. A phenomenological interpretation of the J/ψ A-dependence based on the intrinsic charm ansatz was given in Ref. [45]. It was suggested that large x, $c\bar{c}$ production is dominated by an intrinsic component [43,30], and that the $c\bar{c}$ is "freed" from its virtual state by interactions of the light quarks in the projectile hadron with the target nucleus. The high momentum fraction, small transverse size $c\bar{c}$ cluster in the incident hadron passes through the nucleus undeflected, and it can then evolve into charmonium states after transiting the nucleus. The remaining cluster of light quarks in the intrinsic charm Fock state tends to be absorbed on the front surface of the nucleus. Because the interaction of the light quark components with nucleons is expected to be strong, one then expects that nuclear cross sections will be surface dominated, *i.e.* $\alpha \simeq 2/3$, as seen in the data [46,19,11]. It should be emphasized that this nuclear suppression is unrelated to the shadowing of parton distributions discussed in section 3.

The intrinsic charm model justifies the analysis of Badier *et al.* [19] in which the perturbative and non–perturbative charm production mechanisms were separated on the basis of their different A–dependence ($\alpha = 0.97$ and $\alpha = 0.77$ for a pion beam, respectively). The effective x_F–dependence of α seen in charm production is explained by the different characteristics of the two production mechanisms. Hard gluon fusion production dominates at small x_F, due to the steeply falling gluon structure function. The contribution from intrinsic charm Fock states peaks at higher x_F, due to the large momentum carried by the charm quarks. This two–component hard–scattering plus intrinsic charm model also explains why the nuclear dependence of J/ψ production depends on x_F rather than x_2, as predicted by leading twist factorization [37]. Another important consequence of this picture is that all final states produced by a penetrating intrinsic $c\bar{c}$ component will have the same A–dependence. In particular, the $\psi(2S)$ radially excited state will behave in the same way as the J/ψ, in spite of its larger size. This prediction is confirmed by the recent E–772 data [11]. The nucleus cannot influence the quark hadronization which (at high energies) takes place outside the nuclear environment. In Ref. [7], Vogt, Hoyer, and I have given a systematic analysis of the two-component model including the effects of co-mover interactions.

Recently, Hoyer, Mueller, Tang and I [12] have analysed the role of the various leading and higher twist contributions to heavy quarkonium production in a specific gauge theory model. I will review the main points of this analysis here.

It is important to distinguish three separate kinematical limits for large x heavy quark pairs of invariant mass $\mathcal{M}$: $\mathcal{M} \to \infty$, $x \to 1$ with $\mathcal{M}^2(1-x) \to \infty$; $\mathcal{M} \to \infty$, $x \to 1$ with $\mathcal{M}^2(1-x)$ fixed; and $\mathcal{M} \to \infty$, $x \to 1$ with $\mathcal{M}^2(1-x) \to 0$. In the first limit the usual perturbative QCD scattering formalism applies, including the factorization theorem for hard subprocesses [14]. The effective scale of the interaction, $\mathcal{M}^2(1-x)$, is still asymptotically large. Intrinsic higher twist diagrams, in which the pair is coupled to more than one constituent of the projectile, are damped by powers of $\mathcal{M}^2(1-x)$. In the last limit the standard spectator counting rules for the power behavior of structure functions in the $x \to 1$ limit become valid [47].

It is most interesting to consider the behavior of the QCD processes in the intermediate case in which $\mathcal{M}^2(1-x)$ is held fixed. It turns out that subprocesses involving spectator constituents, which would give power–suppressed higher twist contributions in the first limit, contribute at leading order in the second limit above. Hence the distinction between "extrinsic" processes, where the pair is created by a single gluon, and "intrinsic" processes, where the pair is created by several gluons, essentially disappears. An immediate consequence of the fact that several partons are involved in the leading subprocesses is that the QCD factorization theorem breaks down: Scattering cross sections can no longer be expressed in terms of single parton distributions of the colliding hadrons.

These observations provide a framework for understanding the puzzling phenomena observed at large x, and in particular for explaining the A-dependence of J/ψ production. We find that one can indeed free a virtual $c\bar{c}$ pair, or a lepton pair, at large x by a relatively soft interaction with a light quark component of the projectile. The hardness of the interaction has scale $\mathcal{M}^2(1-x)$, where $\mathcal{M}$ is the pair mass, and the 'freeing' probability is proportional to $1/\mathcal{M}^2(1-x)$. Thus, at sufficiently large x, the cross section for freeing the pair will become large enough so that an x-dependent departure from an A^1 dependence can be expected in the nuclear cross section. If the reaction freeing the pair occurs for an extrinsic, leading-twist component of the projectile, a component of the infinite momentum wavefunction scaling logarithmically with $\mathcal{M}^2$, then the resulting cross section in the above limit is just the normal factorized expression for $q\bar{q}$ annihilation in perturbative QCD. If the freeing occurs for an intrinsic component it corresponds to a higher twist effect in the usual hard scattering formalism.

There are a number of novel features of QCD which emerge in the limit $\mathcal{M} \to \infty$, $x \to 1$ with $\mathcal{M}^2(1-x)$ fixed. In this new QCD limit:

- The leading contributions to the production cross section actually come from spectator interactions rather than direct interactions with the pair itself;
- The coherence of the Fock state is easily broken by soft interactions of finite transverse momentum since the transverse velocity inflicted to the spectators $v_\perp = p_\perp/p(1-x)$ is large;
- QCD factorization is invalid in this limit since there is no relative suppression of interactions involving several constituents of the same hadron;
- The intrinsic mechanism offers the possibility to produce the J/ψ directly in a color singlet state since the three quarks of a proton may all couple to the J/ψ via gluon exchange. Hence there is no need for further soft gluon radiation [48].

- Because of the rapid transverse size expansion of the spectators, production cross sections in nuclear targets become surface dominated at large x;
- The change in physics from the volume-dominated leading-twist fusion subprocesses to surface-dominated higher-twist intrinsic charm contributions occurs as the fractional momentum x of the pair increases. Thus the nuclear target dependence of the production cross section in the fixed $(1-x)\mathcal{M}^2$ limit is a function $A^{\alpha(x)}$ of the pair momentum fraction x rather than a function $A^{\alpha(x_2)}$ of the target parton momentum fraction x_2;
- The change in physics with Feynman x may account for the observed dominance of longitudinally polarized J/ψ at large x.
- Although the absolute normalization of the intrinsic contributions to massive pair production requires knowledge of two or more particle correlations in the non-perturbative hadron wavefunction, it is still possible to use QCD perturbation theory to analyze both the power law behavior and logarithmic evolution of these contributions in a manner analogous to PQCD treatments of large momentum transfer exclusive reactions [32], since the short-distance components of the wavefunction dominate. In particular, the extrapolation from charm to beauty processes is straightforward using scaling at fixed $\mathcal{M}^2(1-x)$.

Our investigation can provide a QCD framework for understanding a number of puzzling features of the large x data:

- The nuclear suppression seen for J/ψ production appears to be larger than that for $\mu^+\mu^-$ production [11]. Only extrinsic diagrams contribute to the lepton pair process, whereas extrinsic and intrinsic contributions contribute to quark pair production.
- Another consequence of large intrinsic contributions is that the charm and beauty structure functions measured in deep inelastic scattering can have a larger than expected support at high x_{Bj}. There is some evidence for this from EMC measurements of the charm structure function of the nucleon [40]. This will be an important area of investigation at HERA. It also implies that charm production near threshold at RHIC could be larger than expected from traditional estimates.
- If the intrinsic charm quarks coalesce with spectators of the projectile, then one can also account for the rather hard momentum distribution of open production observed at large x. Unlike recombination and string pictures, where fast valence quarks pull the heavy particles to high momenta, the intrinsic heavy quark picture implies not only fast open charm and beauty states, but also heavy quarkonium production at large x.
- The intrinsic production mechanism offers an intriguing new possibility of understanding cumulative meson production in nuclei at large transverse momentum [6,5]. As has been established in many experiments, a particle produced in the nuclear fragmentation region can carry more momentum than single nucleons ($x > 1$). The two (or more) gluons transferring momentum to the heavy quarks in intrinsic processes need not, in fact, originate from the same nucleon. Two or more nucleons with a small transverse separation can both transfer momentum to the same quark pair. Moreover, the number of intrinsic diagrams increases quickly with the nuclear number A and the number of gluons involved.
- Among the puzzles of particle production at high x is the observed dependence on the transverse polarization of the beam. Recent data [49] show a remarkable increase of the polarization asymmetry in the region $0.3 < x < 0.9$. It has been difficult to describe such polarization effects in the framework of perturbative QCD because the polarization is given by an interference term between flip and non-flip amplitudes, $Pd\sigma/dt \sim Im(A_f^* A_{nf})$. In leading twist QCD calculations very small

transverse polarizations are expected, since the helicity flip amplitudes are suppressed and the amplitudes are predominantly real. On the other hand intrinsic diagrams have a sizeable imaginary part, which appears immediately in the lowest order amplitude. A similar explanation based on quark-gluon correlations has been given in Ref. [50]. Furthermore, the incoming Fock state involves large transverse momenta for the constituents, which can lead to important helicity flip contributions. It would thus appear that, assuming helicity flip amplitudes are present, the intrinsic production mechanism may be a dominant source of transverse polarization effects. This is supported by the fact that the observed polarization increases with the transverse momentum of the detected hadron. It is in the combined limit of large x and high pair mass (*i.e.* high $p_\perp$) that the intrinsic diagrams contribute at leading order.

- One of the most striking consequences of the analysis of Ref. [12] is the prediction for the soft, low $Q^2 \sim \mathcal{M}^2(1-x)$, Coulomb excitation of the proton to high x massive quark pair configurations, especially $ep \rightarrow e' + J/\psi + X$ and other heavy quarkonium states. The dominant diagrams at $x \rightarrow 1$ involve the electron scattering on the valence spectators, not on the heavy quarks themselves. The absolute size of the cross section can be estimated from J/ψ production in proton–proton collisions, assuming factorization. A critical test of the importance of higher twist correlations in the proton wavefunction will be the observation of quarkonium states at large x in the proton's fragmentation region. Furthermore, if intrinsic diagrams account for the transverse polarization observed in high x hadronic reactions, then similar polarization effects should be present universally in the target fragmentation region whether the reaction is peripheral ep collisions or diffractive heavy ion collisions.

6. The Nuclear Fragmentation Region

It is clear that many of the most interesting tests of QCD in relativistic heavy ion collisions will require measurements of particle production over the full rapidity range. Studies of massive pair and heavy quark production at large x, color transparency tests, cumulative reactions, diffractive jet production, and coherent nuclear reactions which test the nature of the Pomeron coupling to nuclei [51] all require observation of produced hadrons or the recoil nucleus in the target and projectile fragmentation regions.

We have also noted that the physics of the central rapidity region in high energy nucleus-nucleus collisions will be severely complicated by minijet production and hadronization. The interactions between minijets and the interactions of minijets with the remaining spectators is certainly of interest, including possible non-linear hadronization effects due to the scattering and coalescence with co-movers. However, the clearest searches for a change of phase from nuclear or hadronic matter to quark-gluon degrees of freedom may well involve the nuclear fragmentation regions where the contributions due to minijet production are less important.

The physics of the target fragmentation region is also interesting from the standpoint of light-cone physics. As one nucleus passes through the other in a relativistic heavy ion collision, the structure of the target nucleus is excited at a fixed "light-cone" time $\tau = t + z/c$ rather than a fixed ordinary time t. Thus the boundary condition for "heating" the nucleus is set by light-cone time τ. The evolution operator in light-cone time τ is the light-cone Hamiltonian $P^- = P^0 - P^z$ which has eigenvalues $(\mathcal{M}^2 + P_\perp^2)/(P^0 + P^z)$, where $\mathcal{M}$ is the invariant mass spectrum of the system. It thus should be advantageous to study statistical and thermodynamic quantities such as the covariant partition function $Z_{cov} = \sum_{\mathcal{M}} exp-(\mathcal{M}^2/T^*)$, where T^* plays the role of a frame-independent temperature.

In the case of simple theories such as QCD in one-space and one-time, Z_{cov} can be evaluated explicitly since the complete spectrum can be obtained using matrix diagonalization methods [52].

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