HIGH DENSITY QCD, SATURATION AND DIFFRACTIVE DIS

I.P. IVANOV  
Institute of Mathematics, Novosibirsk, Russia  

N.N. NIKOLAEV\textsuperscript{A,B), W. SCHÄFER\textsuperscript{A)}, B.G. ZAKHAROV\textsuperscript{B)}  
\textsuperscript{A)} Institut f. Kernphysik, Forschungszentrum Jülich,  
D-52425 Jülich, Germany  
\textsuperscript{B)} L.D. Landau Institute for Theoretical Physics,  
Chernogolovka, Russia  

AND  
V.R. ZOLLER  
Institute for Theoretical and Experimental Physics,  
Moscow, Russia

Abstract

We review a consistent description of the fusion and saturation of partons in the Lorentz-contracted ultrarelativistic nuclei in terms of a nuclear attenuation of color dipole states of the photon and collective Weizsäcker-Williams (WW) gluon structure function of a nucleus. Diffractive DIS provides a basis for the definition of the WW nuclear glue. The point that all observables for DIS off nuclei are uniquely calculable in terms of the nuclear WW glue amounts to a new form of factorization in the saturation regime.

1. Introduction

Within the QCD parton model the virtual photoabsorption cross section is proportional to the density of partons in the target and vice versa. When DIS is viewed in the laboratory frame, the hadronic properties of photons suggest [1] a nuclear shadowing and depletion of the density of partons, when DIS is viewed in the Breit frame, the Lorentz contraction of an ultrarelativistic nucleus entails a spatial overlap and fusion of partons at $x \lesssim x_A = 1/R_A m_N \sim 0.1 \cdot A^{-1/3}$. This interpretation of nuclear opacity in
terms of a fusion and saturation of nuclear partons has been introduced in 1975 [1] way before the QCD parton model. The pQCD link between nuclear opacity and saturation has been considered in ref. [2] and by Mueller [3], the pQCD discussion of fusion of nuclear gluons has been revived by McLerran et al. [4].

Amplitudes of diffractive DIS are intimately related to the unintegrated glue of the target [5, 6]. Because coherent diffractive DIS in which the target nucleus does not break makes precisely 50 per cent of the total DIS events for heavy nuclei at small $x$ [7], diffractive DIS off nuclei offers a unique definition of the collective WW glue of Lorentz-contracted ultrarelativistic nuclei [8]. The recent work has shown that all the observables for nuclear DIS are uniquely calculable in terms of the NSS-defined WW nuclear glue [9, 10]. In this overview presented at Diffraction'2002 by one of the authors (N.N.N) we summarize this new development.

2. Quark and antiquark jets in DIS off free nucleons: single particle spectrum and jet-jet decorrelation

In the color dipole approach to DIS [2, 5, 7, 11, 12, 13] the fundamental quantity is the cross section for interaction of the $q\bar{q}$ dipole on a nucleon,

$$\sigma(r) = \alpha_S(r)\sigma_0 \int d^2\kappa f(\kappa) [1 - \exp(i\kappa r)]$$ (1)

where $f(\kappa)$ is related to the unintegrated glue of the target nucleon by

$$f(\kappa) = \frac{4\pi}{N_c\sigma_0} \cdot \frac{1}{\kappa^4} \cdot \frac{\partial G}{\partial \log \kappa^2}.$$ (2)

For DIS off a free nucleon, see figs. 1a-1d, $\sigma_N = \int d^2r d^2z |\Psi(z,r)|^2 \sigma(r)$, and the jet-jet inclusive cross section equals

$$\frac{d\sigma_N}{dzd^2p_+d^2\Delta} = \frac{\sigma_0}{2} \cdot \frac{\alpha_S(p_+^2)}{(2\pi)^2} f(\Delta) \left| \langle \gamma^* | p_+ \rangle - \langle \gamma^* | p_+ - \Delta \rangle \right|^2$$ (3)

where $p_+$ is the transverse momentum of the quark, $\Delta = p_+ + p_-$ is the jet-jet decorrelation momentum, $z_+ = z$ and $z_- = 1 - z$ are the fractions of photon’s lightcone momentum carried by the quark and antiquark, respectively, and the photon wave functions $\Psi(r)$ and $\langle p|\gamma^* \rangle$ are found in [2, 5]. The crucial point is that the jet-jet decorrelation is controlled [14] by the unintegrated gluon SF $f(\Delta)$.

3. Non-Abelian propagation of color dipoles in nuclear medium

DIS at $x \lesssim x_A$ is dominated by interactions of $q\bar{q}$ states of the photon. The unitarity cuts of the free-nucleon forward Compton diagrams of figs.
Figure 1. The pQCD diagrams for inclusive (a-d) and diffractive (e,f) DIS off protons and nuclei (g-k). Diagrams (a-d) show the unitarity cuts with color excitation of the target nucleon, (g) - a generic multiple scattering diagram for Compton scattering off nucleus, (h) - the unitarity cut for a coherent diffractive DIS, (i) - the unitarity cut for quasielastic diffractive DIS with excitation of the nucleus $A^*$, (j,k) - the unitarity cuts for truly inelastic DIS with single and multiple color excitation of nucleons of the nucleus.

1a-1d describe excitation from the color-neutral to color-octet $q\bar{q}$ pair. The unitarity cuts of the nuclear Compton scattering amplitude of fig. 1g which correspond to diffractive and genuine inelastic DIS with color excitation of the nucleus are shown in fig. 1. The elegant multichannel description of the non-Abelian intranuclear evolution of color dipoles is found in [10], here we cite only the principal results.

Let $b_+$ and $b_-$ be the impact parameters of the quark and antiquark, respectively, and $S_A(b_+, b_-)$ be the S-matrix for interaction of the $q\bar{q}$ pair with the nucleus. We are interested in the 2-body (jet-jet) inclusive inelastic cross section when we sum over all color excitations of the target nucleus and all color states $c_{km}$ of the $q_k\bar{q}_m$ pair:

$$\frac{d\sigma_{im}}{dzdp_+dp_-} = \frac{1}{(2\pi)^3} \int d^2b'_+d^2b'_-d^2b_+d^2b_- \times \exp[-ip_+(b_+ - b'_+) - ip_-(b_- - b'_-)]\Psi^*(b'_+ - b'_-)\Psi(b_+ - b_-)$$
\[
\times \left\{ \sum_{A^*} \sum_{km} (1; A|S_A^* (b'_+, b'_-)|A^*; c_{km}) \langle c_{km}; A^*|S_A (b_+, b_-)|A; 1 \rangle \\
- \langle 1; A|S_A^* (b'_+, b'_-)|A; 1 \rangle \langle 1; A|S_A (b_+, b_-)|A; 1 \rangle \right\},
\]

where the diffractive component of the final state has been subtracted. According to [9, 10], upon summing over nuclear final states \( A^* \) and making use of the technique developed in [15, 16], the integrand of (4) can be represented as an \( S \)-matrix \( S_{4A} (b_+, b_-, b'_+, b'_-) \) for the propagation of the two quark-antiquark pairs in the overall singlet state. The detailed description of the matrix of 4-parton color dipole cross section \( \sigma_4 \) is found in [10]. The single-jet cross section evaluated at the parton level equals [9]

\[
\frac{d\sigma_{in}}{d^2 b d^2 p dz} = \frac{1}{(2\pi)^2} \int d^2 r d^2 r \exp[i p(r' - r)] \Psi^*(r') \Psi(r) \\
\times \left\{ \exp[-\frac{1}{2} \sigma(r - r') T(b)] - \exp[-\frac{1}{2} [\sigma(r) + \sigma(r')] T(b)] \right\},
\]

(5)

4. The Pomeron-Splitting Mechanism for Diffractive Hard Dijets and Weizsäcker-Williams glue of nuclei

The two distinct diffractive dijet production QCD subprocesses are the classic Landau-Pomeranchuk-Feinberg-Glauber beam-splitting [17] (fig. 1e, fig. 2a) and the Nikolaev-Zakharov pomeron-splitting [5, 6] (fig. 1f, fig. 2b):

\[
\Phi_0(z, p) = \int d^2 r e^{-i p r} \sigma(r) \Psi(z, r) \\
= \alpha_S(p^2) \sigma_0 \left[ \langle p|\gamma^* \rangle \int d^2 \kappa f(\kappa) - \int d^2 \kappa \langle \kappa|\gamma^* \rangle f(p - \kappa) \right],
\]

(6)

The amplitude for the former mechanism is \( \propto \langle p|\gamma^* \rangle \) and the transverse momentum \( p \) of jets comes from the intrinsic transverse momentum of \( q, \bar{q} \) in the beam particle, in the latter jets receive a transverse momentum from gluons in the Pomeron, notice the convolution of the wave function with the unintegrated glue in the target proton. If the beam particle were a pion, then \( \psi(z, p) \) would be much steeper than \( f(p) \), and the asymptotics of the convolution integral will be [8]

\[
\int d^2 \kappa \psi_\pi(z, \kappa) f(p - \kappa) \approx f(p) \int d^2 \kappa \psi_\pi(z, \kappa) = f(p) \phi_\pi(z) F_\pi,
\]

(7)

and, furthermore, will probe the pion distribution amplitude \( \phi_\pi(z) \).
Now we notice that in the color dipole representation the nuclear amplitude is readily obtained [2] by substituting in eq. (6)

$$\sigma(r) \rightarrow \sigma_A(r) = 2 \int d^2b \{1 - \exp[-\nu_A(b)]\}$$

(8)

$$\nu_A(b) = \frac{1}{2} \alpha_S(r) \sigma_0 T(b) = \frac{1}{2} \alpha_S(r) \sigma_0 \int dz n_A(b, z)$$

(9)

Typical nuclear double scattering diagrams of figs. 2c-2e can be classified as shadowing of the pion splitting (fig. 2c), shadowing of single Pomeron splitting (fig. 2d) and double Pomeron splitting (fig. 2e) contributions. In the $j$ Pomeron splitting $j$ exchanged Pomerons couple with one gluon to the quark and with one gluon to the antiquark of the dipole. That involves the $j$-fold convolution $f^{(j)}(\kappa) = \prod_{i=1}^{j} d^2\kappa_i f(\kappa_i) \delta(\kappa - \sum_{i=1}^{j} \kappa_i)$. Now we can invoke the NSS representation [8, 9]

$$\exp[-\nu(b)] = \int d^2\kappa \Phi(\nu_A(b), \kappa) \exp(i\kappa s)$$

(10)

where $f^{(0)}(\kappa) = \delta(\kappa)$, $\Phi(\nu_A(b), \kappa) = \exp(-\nu_A(b)) f^{(0)}(\kappa) + \phi_{WW}(b, \kappa)$ and

$$\phi_{WW}(b, \kappa) = \exp[-\nu_A(b)] \sum_{j=1}^{\infty} \frac{1}{j!} \nu_A^{(j)}(b) f^{(j)}(\kappa)$$

(11)

can be identified with the unintegrated nuclear Weizsäcker-Williams glue per unit area in the impact parameter plane [9].
5. Nuclear dilution and broadening of the unintegrated Weizsäcker-Williams glue of nuclei

The hard tail of WW glue per bound nucleon is calculable parameter free:

$$f_{WW}(b, \kappa) = \frac{\phi_{WW}(b, \kappa)}{\nu_A(b)} = f(\kappa) \left[ 1 + \frac{2C_A\pi^2\gamma^2\alpha_S(r)T(b)}{C_FN_c\kappa^2}G(\kappa^2) \right]$$ \hspace{1cm} (12)

In the hard regime the differential nuclear glue is not shadowed, furthermore, because of the manifestly positive-valued and model-independent nuclear higher twist correction it exhibits nuclear antishadowing property [8]. The application of this formalism to the interpretation of the E791 data on coherent diffraction of pions into dijets [18] is found in [8, 19].

![Figure 3. The dilution for soft momenta and broadening for hard momenta of the multiple convolutions $f^{(j)}(k)$.](image)

In the soft region one can argue that

$$\phi_{WW}(\kappa) \approx \frac{1}{\pi} \frac{Q_A^2}{(\kappa^2 + Q_A^2)^2},$$ \hspace{1cm} (13)

where the saturation scale $Q_A^2 = \nu_A(b)Q_0^2 \propto A^{1/3}$. Notice a strong nuclear dilution of soft WW glue, $\phi_{WW}(\kappa) \propto 1/Q_A^2 \propto A^{-1/3}$, which must be contrasted to the $A$-dependence of hard WW glue (12),

$$\phi_{WW}(\kappa) = \nu_A(b)f_{WW}(\kappa) \propto A^{1/3} \times (1 + A^{1/3} \times (HT) + ...).$$ \hspace{1cm} (14)

Take the platinum target, $A = 192$. The numerical estimates based on the parameterization [20] for $f(\kappa)$ show that for $(q\bar{q})$ color dipoles in the
average DIS on this target $\nu_A \approx 4$ and $Q_{2A}^2 \approx 0.8\,\text{GeV}^2$. For the $q\bar{q}g$ Fock states of the photon, which behave predominantly like the dipole made of the two octet color charges, $Q_{2A}^2$ is larger by the factor $C_A/C_F = 9/4$, and for the the average DIS on the Pt target $Q_{2A}^2 \approx 2.2\,\text{GeV}^2$.

6. Nuclear partons in the saturation regime

On the one hand, making use of the NSS representation, the total nuclear photoabsorption cross section (9) can be cast in the form

$$\sigma_A = \int d^2b \int dz \int \frac{d^2p}{(2\pi)^2} \int d^2\kappa \phi_{WW}(\kappa) \langle \gamma^*|p\rangle - \langle \gamma^*|p-\kappa\rangle|^2$$

(15)

and its differential form defines the IS sea in a nucleus,

$$\frac{d\bar{q}_{IS}}{d^2bd^2p} = \frac{1}{2} \frac{Q^2}{4\pi^2\alpha_{em}} \frac{d\sigma_A}{d^2bd^2p}.$$  

(16)

Remarkably, in terms of the WW nuclear glue, all intranuclear multiple-scattering diagrams of fig. 3g sum up to precisely the same four diagrams fig. 3a-3d as in DIS off free nucleons. On the other hand, making use of the NSS representation, after some algebra one finds [9]

$$\frac{d\sigma_{in}}{d^2bd^2pdz} = \frac{1}{(2\pi)^2} \left\{ \int d^2\kappa \phi_{WW}(\kappa) \langle \gamma^*|p\rangle - \langle \gamma^*|p-\kappa\rangle|^2 \right. \\
- \left. \int d^2\kappa \phi_{WW}(\kappa) \left( \langle \gamma^*|p\rangle - \langle \gamma^*|p-\kappa\rangle \right) \right\}.$$  

(17)

$$\frac{d\sigma_D}{d^2bd^2pdz} = \frac{1}{(2\pi)^2} \left| \int d^2\kappa \phi_{WW}(\kappa) \langle \gamma^*|p\rangle - \langle \gamma^*|p-\kappa\rangle \right|^2.$$  

(18)

Putting the inelastic and diffractive components of the FS quark spectrum together, one finds the FS parton density which exactly coincides with the IS parton density (16)!

Consider first $p^2 \lesssim Q^2 \lesssim Q_{2A}^2$ when the nucleus is opaque for all color dipoles in the photon. In this regime the nuclear counterparts of diagrams of figs. 1b,1d,1f can be neglected and diffraction will be dominated by the the Landau-Pomeranchuk mechanism of fig. 1e,2a:

$$\frac{d\bar{q}_{FS}}{d^2bd^2p} \approx \frac{Q^2}{8\pi^2\alpha_{em}} \int dz \left| \int d^2\kappa \phi_{WW}(\kappa) \right|^2 |\langle \gamma^*|p\rangle|^2 \approx \frac{N_c}{4\pi^4}. $$

(19)

Remarkably, diffractive DIS measures the momentum distribution in the $q\bar{q}$ Fock state of the photon. In contrast to diffraction off free nucleons [5, 6, 21],
diffraction off opaque nuclei is dominated by the anti-collinear splitting of hard gluons into soft sea quarks, \( \kappa^2 \gg p^2 \). Precisely for this reason one finds the saturated FS quark density, because the nuclear dilution of the WW glue is compensated for by the expanding plateau.

The related analysis of the FS quark density for truly inelastic DIS in the same domain of \( p^2 \lesssim Q^2 \lesssim Q_A^2 \) gives

\[
\frac{dq_{FS}}{d^2b d^2p_{bm}} = \frac{1}{2} \frac{Q^2}{4\pi^2\alpha_{em}} \int dz \int d^2\kappa \phi_{WW}(\kappa) |\langle \gamma^* | p - \kappa \rangle|^2 \\
\approx \frac{Q^2}{8\pi^2\alpha_{em}} \phi_{WW}(0) \int d^2\kappa \int dz |\langle \gamma^* | \kappa \rangle|^2 \\
= \frac{N_c}{4\pi^4} \frac{Q^2}{Q_A^2} \cdot \theta(Q_A^2 - p^2)
\]

which, as a functional of the photon wave function and nuclear WW gluon distribution, is completely different from the free-nucleon version of (15).

Figure 4. The two-plateau structure of the momentum distribution of FS quarks: (i) \( Q^2 \lesssim Q_A^2 \): The inelastic plateau is much broader than the diffractive one. (ii) \( Q^2 \gtrsim Q_A^2 \): inelastic plateau is identical to diffractive one.

Now notice, that in the opacity regime the diffractive FS parton density coincides with the contribution \( \propto |\langle \gamma^* | p \rangle|^2 \) to the IS sea parton density from the spectator diagram 1a, whereas the FS parton density for truly inelastic DIS coincides with the contribution to IS sea partons from the diagram of fig. 1c. The contribution from the crossing diagrams 1b,d is
negligibly small. Nuclear broadening and unusually strong $Q^2$ dependence of the FS/IS parton density from truly inelastic DIS, demonstrate clearly a distinction between diffractive and inelastic DIS, see fig. 4.

In the case of soft quarks, $p^2 \lesssim Q^2_A$, in hard photons, $Q^2 \gg Q^2_A$, the result (19) for diffractive DIS is retained, whereas in the numerator of the result (20) for truly inelastic DIS one must substitute $Q^2 \rightarrow Q^2_A$, so that in this case $d\bar{q}_{FS\mid D} \approx d\bar{q}_{FS\mid in}$ and $d\bar{q}_{IS} \approx 2d\bar{q}_{FS\mid D}$, see fig. 4. The evolution of soft nuclear sea, $p^2 \lesssim Q^2_A$, is entirely driven by an anti-collinear splitting of the NSS-defined WW nuclear glue into the sea partons.

The early discussion of the FS quark density in the saturation regime is due to Mueller [22]. Mueller focused on $Q^2 \gg Q^2_A$ and discussed neither a distinction between diffractive and truly inelastic DIS nor a $Q^2$ dependence and broadening (20) for truly inelastic DIS at $Q^2 \sim Q^2_A$.

7. Signals of saturation in exclusive diffractive DIS

The flat $p^2$ distribution of forward $q, \bar{q}$ jets in truly inelastic DIS in the saturation regime must be contrasted to the $\propto G(p^2)/p^2$ spectrum for the free nucleon target. In the diffractive DIS the saturation gain is much more dramatic: flat $p^2$ distribution of forward $q, \bar{q}$ jets in diffractive DIS in the saturation regime must be contrasted to the $\propto 1/(p^2)^2$ spectrum for the free nucleon target [5, 6, 21]. In the general case one must compare the saturation scale $Q_A$ to the relevant hard scale $\bar{Q}^2$ for the specific diffractive process. For instance, in exclusive diffractive DIS, i.e., the vector meson production, $\bar{Q}^2 \approx (Q^2 + m_V^2)/4$ and the transverse cross section has been predicted to behave as [23]

$$\sigma_T \propto G^2(x, \bar{Q}^2)(\bar{Q}^2)^{-4} \quad (21)$$

At $\bar{Q}^2 > Q^2_A$ the same would hold for nuclei too, but in the opposite case of $\bar{Q}^2 < Q^2_A$ the $\bar{Q}^2$-dependence is predicted to change to

$$\sigma_T \propto G^2(x, \bar{Q}^2)(Q^2_A)^{-2}(\bar{Q}^2)^{-2} \quad (22)$$

8. Jet-jet decorrelation in DIS off nuclear targets

The derivation of the jet-jet inclusive cross section requires a full fledged non-Abelian multichannel calculation of $S_{1A}(b_+, b_-, b'_+, b'_-)$. The principal result for the hard, $p^2_\pm \gg Q^2_A$, jet-jet cross section is [10]

$$\frac{d\sigma_{in}}{d^2 b dz d^2 p_+ d^2 \Delta} = T(b) \int_0^1 d\beta \int d^2 \kappa \Phi(2\beta \lambda_\nu A(b), \Delta - \kappa) \frac{d\sigma_N}{dz d^2 p_+ d^2 \kappa} \quad (23)$$
It has a probabilistic form of a convolution of the differential cross section on a free nucleon target with the unintegrated nuclear WW glue $\Phi(2\beta \lambda c\nu_A(b), \kappa)$, $\beta$ has a meaning of the fraction of the nuclear thickness which the ($q\bar{q}$) pair propagates in the color-octet state and the nontrivial color factor $\lambda_c = (N_c^2 - 2)/(N_c^2 - 1)$ is a manifestation of the non-Abelian intranuclear propagation of color dipoles.

The azimuthal decorrelation of hard jets is quantified by

$$\langle \Delta^2 \rangle \approx \frac{Q_A^2}{2} \left( \frac{Q_A^2 + p^2_\perp}{p^2_\perp} \log \frac{Q_A^2 + p^2_\perp}{Q_A^2} - 1 \right)$$

and the differential out-of-plane momentum distribution

$$\frac{d\sigma}{d\Delta^\perp} \approx \frac{1}{2} \frac{Q_A^2}{(Q_A^2 + \Delta^\perp)^{3/2}}$$

which shows that the probability to observe the away jet at $\Delta^\perp \sim 0$ disappears $\propto 1/Q_A \propto 1/\sqrt{\nu_A}$. For instance, from peripheral DIS to central DIS on a heavy nucleus like $Pt$, $\nu_{8A}$ rises from $\nu_{8A} = 1$ to $\nu_{8A} \sim 13$, so that according to (25) a probability to find the away jet decreases by the large factor $\sim 3.5$ from peripheral to central DIS off $Pt$ target.

In the generic case the closed form for the jet-jet inclusive cross section is found [10] only for large $N_c$:

$$\frac{d\sigma_{in}}{d^2bzd\Delta} = \frac{1}{2(2\pi)^2} \alpha_S \sigma_0 T(b) \int_0^1 d\beta \int d^2\kappa_1 d^2\kappa_2 d^2\kappa_3 d^2\kappa f(\kappa) \times \Phi((1 - \beta)\nu_A(b), \kappa_1) \Phi((1 - \beta)\nu_A(b), \kappa_2) \times \Phi(\beta\nu_A(b), \kappa_3) \Phi(\beta\nu_A(b), \Delta - \kappa_3 - \kappa) \times \{\Psi(-p_\perp + \kappa_2 + \kappa_3) - \Psi(-p_\perp + \kappa_2 + \kappa_3 + \kappa)\}^* \times \{\Psi(-p_\perp + \kappa_1 + \kappa_3) - \Psi(-p_\perp + \kappa_1 + \kappa_3 + \kappa)\}$$

We emphasize that it is uniquely calculable in terms of the NSS-defined WW glue of the nucleus.

One important implication of (26) is that for minijets with the jet momentum $|p_\perp|, |\Delta| \lesssim Q_A$, the the minijet-minijet inclusive cross section depends on neither the minijet nor decorrelation momentum, i.e., we predict a complete disappearance of the azimuthal decorrelation of jets with the transverse momentum below the saturation scale.

Our experience with application of color dipole formalism to hard hadron-nucleus interactions [15, 16] suggests that our analysis can be readily generalized to mid-rapidity jets. For this reason, we expect similar strong decorrelation of mid-rapidity jets in hadron-nucleus and nucleus-nucleus collisions. To this end, recently the STAR collaboration reported a disappearance
of back-to-back high $p_\perp$ hadron correlation in central gold-gold collisions at RHIC [24]. Nuclear enhancement of the azimuthal decorrelation of the trigger and away jets may contribute substantially to the STAR effect.

9. Summary and conclusions

We reviewed a theory of DIS off nuclear targets based on the consistent treatment of propagation of color dipoles in nuclear medium. What is viewed as attenuation in the laboratory frame can be interpreted as a fusion of partons from different nucleons of the ultrarelativistic nucleus. Diffractive attenuation of color single $q\bar{q}$ states gives a consistent definition of the WW unintegrated gluon structure function of the nucleus [8, 9], all other nuclear DIS observables - sea quark structure function and its decomposition into equally important genuine inelastic and diffractive components, exclusive diffraction off nuclei, the jet-jet inclusive cross section, - are uniquely calculable in terms of the NSS-defined nuclear WW glue. This property can be considered as a new factorization which connects DIS in the regimes of low and high density of partons. The anti-collinear splitting of WW nuclear glue is a clearcut evidence for inapplicability of the DGLAP evolution to nuclear structure functions unless $Q^2 \gg Q^2_{8A}$.

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