Decisive test of color transparency in exclusive electroproduction of vector mesons

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Abstract

The exclusive production of vector mesons in deep inelastic scattering is a hard scattering process with the well controlled size of quark configurations which dominate the production amplitude. This allows an unambiguous prediction of color transparency effects in the coherent and incoherent production of vector mesons on nuclei. We demonstrate how the very mechanism of color transparency leads to a belated onset of color transparency effects as a function of $Q^2$. We conclude that the $Q^2$ dependence of the exclusive $\rho^0$-meson production on nuclei and nucleons observed in the Fermilab E665 experiment gives a solid evidence for the onset of color transparency. We propose the scaling relation between the $\rho^0$ and the $J/\Psi$ production, which further tests the mechanism of color transparency in exclusive (virtual) photoproduction.

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The exclusive production of vector mesons $\gamma^* p \to V p$ in deep inelastic scattering is a hard scattering process in which the transverse size $r_Q$ of quark configurations which dominate the production amplitude is under the good control \[1\],

$$r_Q = \frac{2}{\sqrt{m_V^2 + Q^2}}, \quad (1)$$

which makes this reaction an ideal laboratory for testing color transparency (CT) ideas (for the recent review on CT see \[2,3\]). In the exclusive production of nuclei, the nuclear attenuation effects were predicted \[1\] to be $\propto r_Q^2$ and vanish at large virtuality $Q^2$ of the photon, which was confirmed by the recent data on the $\rho^0$ production from the FNAL E665 experiment \[4\]. In \[1\] we have presented quantitative predictions for the incoherent production of $\rho^0$-mesons on iron nucleus. Since the E665 data give the first definitive signal of CT in the hard scattering process, an analysis of the further consequences of CT for virtual photoproduction of vector mesons is called upon. In this paper we present predictions for the signal of CT in the coherent production on nuclei and extend calculations \[1\] for the incoherent production to wider range of nuclei, using the formalism developed by us earlier \[1,5,6\]. We demonstrate how the very mechanism of CT leads to a belated onset of CT effects as a function of $Q^2$, in good agreement with the E665 data \[4\]. We propose the scaling law which relates CT effects in the (virtual) photoproduction of the $\rho^0$ and the $J/\psi$ mesons.

We begin with a brief outline of the lightcone approach to exclusive $\gamma^* N \to VN$ production developed in \[5,6,1,2\]. At the high energy $\nu$ of the virtual photon the reaction mechanism greatly simplifies: The photon fluctuates into the $q\bar{q}$ pair at a large distance (the coherence length)

$$l_c = \frac{2\nu}{Q^2 + m_V^2} \quad (2)$$

in front of the target nucleon (nucleus). After interaction the $q\bar{q}$ pair recombines into the vector meson $V$ with the recombination (formation) length

$$l_f = \frac{\nu}{m_V \Delta m} \quad (3)$$

where $\Delta m$ is the typical level splitting in the quarkonium. At high energy $\nu$ both $l_c$ and $l_f$ greatly exceed the radius $R_N$ ($R_A$) of the target nucleon (nucleus), and the transverse size $\vec{r}$ of the $q\bar{q}$ pair and the longitudinal momentum partition $z$ and $(1 - z)$ between the quark and antiquark of the pair do not change during the interaction with the target. This enables one to introduce the lightcone wave function $\Psi_{\gamma^*}(\rho, z)$ of the $q\bar{q}$ fluctuation \[7\]. The color-singlet $q\bar{q}$ pair interacts with the target nucleon with the cross section \[7,8\]

$$\sigma(r) = \frac{\pi^2}{3} r^2 \mathcal{F}(\nu, r) \quad (4)$$

where $\mathcal{F}(\nu, r)$ is related at small $r$ to the gluon structure function of the proton, $\mathcal{F}(\nu, r) = \alpha_S(r) x g(x, Q_r^2)$, evaluated at $Q_r^2 \sim 1/r^2$ and the value of the Bjorken variable $x \approx (Q_r^2 + m_V^2)/2m_N \nu$, and $\alpha_S(r)$ is the running QCD coupling. This factor $\mathcal{F}(\nu, r)$ takes into account the effect of higher $q\bar{q}g_1...g_n$ Fock components in the lightcone wave function of the photon and the vector meson \[8\]. It is a smooth function of $r$ compared to $r^2$ in Eq. (4).
cross section $\sigma(r)$ has the CT property of vanishing at $r \to 0$ [9,10]. Since the CT property is conveniently quantified in terms of $\sigma(r)$ (for the review see [2,3]), it is important to understand how $\sigma(r)$ is probed in the exclusive electroproduction.

In order to set up the reference frame notice [1,2,5,6], that the amplitude of the forward photoproduction $\gamma^* N \to VN$ can be cast in the quantum-mechanical form ($\vec{q}$ is the momentum transfer)

$$M(VN, \vec{q} = 0) = \langle V|\sigma(r)|\gamma^* \rangle = \int_0^1 dz \int d^2\vec{r}\sigma(r)\Psi_V(r,z)^*\Psi_{\gamma^*}(r,z).$$

(5)

The most important feature of $\Psi_{\gamma^*}(r,z)$ is an exponential decrease at large distances [7],

$$\Psi_{\gamma^*}(r,z) \propto \exp(-\varepsilon r),$$

(6)

where

$$\varepsilon^2 = m_q^2 + z(1 - z)Q^2 \approx \frac{1}{4}(m_V^2 + 4z(1 - z)Q^2).$$

(7)

In the nonrelativistic quarkonium $z \approx 1/2$, $\Psi_{\gamma^*}(r,z)$ is concentrated at $r \lesssim r_Q$ and the wave function (6) [7] and Eq. (5) for $r_Q$ [1] fulfill the dream [11] of having the well specified control of the size of quark configurations important in the hard scattering process.

The generalization to nuclear targets is straightforward [1,5,6]. In the incoherent production $\gamma^* A \to VA^*$ one sums over all excitations of the final nucleus. The nuclear transparency for the incoherent production equals \([12,2,5,6]\)

$$T_A = \frac{\sigma_A}{A\sigma_p} = \frac{1}{A} \int d^2\vec{b} T(b) \frac{\langle V|\sigma(r)\exp\left[-\frac{1}{2}\sigma(r)T(b)\right]|\gamma^* \rangle^2}{\langle V|\sigma(r)|\gamma^* \rangle^2} = 1 - \Sigma_V \frac{1}{A} \int d^2\vec{b} T(b)^2 + \ldots$$

(8)

where $T(b) = \int dz n_A(b,z)$ is the optical thickness of a nucleus at the impact parameter $b$, the nuclear density $n_A(b,z)$ is normalized to the nuclear mass number $A$, $\int d^2\vec{r} n_A(\vec{r}) = A$, (for the compilation of the nuclear density parameterizations see [13]). The observable [2]

$$\Sigma_V = \frac{\langle V|\sigma(r)^2|\gamma^* \rangle}{\langle V|\sigma(r)|\gamma^* \rangle},$$

(9)

measures the strength of intranuclear final state interaction (FSI).

The amplitude of the coherent nuclear production $\gamma^* A \to VA$ equals [9,12]

$$M(VA, \vec{q}) = 2 \int d^2\vec{b} \langle V|1 - \exp\left[-\frac{1}{2}\sigma(r)T(b)\right]|\gamma^* \rangle \exp(-i\vec{q}\vec{b})$$

$$= AM(VN, \vec{q} = 0)[G_{em}(q) - \Sigma_V \frac{G_2(q)}{2A} \int d^2\vec{b} T(b)^2 + \ldots]$$

(10)

Here $G_{em}(q)$ is the charge form factor of the nucleus and the form factor of double scattering $G_2(q)$ is defined by $\int d^2\vec{b} T(b)^2 \exp(-i\vec{q}\vec{b}) = G_2(q) \int d^2\vec{b} T(b)^2$. The total cross section of the coherent production equals

$$\sigma_{coh}(VA) = 4 \int d^2\vec{b} \left|\langle V|1 - \exp\left[-\frac{1}{2}\sigma(r)T(b)\right]|\gamma^* \rangle\right|^2$$

$$= |M(VN, \vec{q} = 0)|^2 \int d^2\vec{b} T(b)^2 \left[1 - \frac{1}{2}\Sigma_V T(b) + \ldots\right]$$

(11)
For the sake of clarity in the subsequent discussion of the onset of CT, in Eqs. (8,11,12) we have explicitly shown the leading term of FSI. The signature of the coherent production is a sharp forward diffraction peak with the slope $B_{coh} \approx \frac{1}{3} R_{ch}(A)^2 \gg B(VN)$, see Eq. (10), whereas in the incoherent nuclear production the $q^2$-dependence is to a good accuracy the same as in production on the free nucleons ([14], for the review see [15]). This allows to separate the coherent and incoherent cross sections even at limited resolution in $q^2$.

Now we turn to implications of Eq. (1) and of CT (4) for the production rate and nuclear FSI effects. The wave function of the vector meson is smooth at small $r$. Because of CT property (4) the integrand of (5) is $\propto r^3 \exp(-\frac{r}{r_Q})$ and the amplitude of (virtual) photoproduction will be dominated by contribution from size $r_S \sim 3r_Q$, (12)

which falls in into the perturbative QCD domain $r_S \ll R_V$ at a sufficiently large $Q^2$. In this domain for production of the transversely polarized vector mesons by the transversely polarized photons $\gamma^*_T N \to V_T N$ we have an estimate

$$M_T(VN, \vec{q} = 0) \propto \frac{r_S^2}{R_V^{3/2}} \sigma(r_S) \propto \frac{1}{(Q^2 + m_V^2)^2} F(\nu, r_S).$$

(13)

The $(m_V^2 + Q^2)^{-2}$ behavior of the amplitude (13) is different from the VDM prediction [15,16]. This difference comes from the factor $\sigma(r_S) \sim 1/(Q^2 + m_V^2)$ in (13) which emphasizes a relevance of CT property of $\sigma(r)$ to the total production rate.

The evaluation of the strength of FSI $\Sigma_\nu$ goes as follows: The integrand of the matrix element $\langle V | \sigma(r^2) | \gamma^* \rangle$ is $\sim r^5 \exp(-\frac{r}{r_Q})$ and is peaked at

$$r \sim r_{FSI} = (4 - 5)r_Q.$$

(14)

(Extension to the higher-order rescatterings is straightforward.) This gives an estimate

$$\Sigma_\nu \approx \sigma(r_{FSI}).$$

(15)

CT and/or weak FSI set in when $\Sigma_\nu \ll \sigma_{tot}(VN) \approx \sigma(R_V)$, i.e., when $r_{FSI} \ll R_V$. Remarkably, the large numerical factor $\approx (4 - 5)$ in the r.h.s. of Eq. (14) comes from CT property of $\sigma(r)$, and the very CT property of the production mechanism predicts a belated onset of the CT effect in nuclear attenuation, which requires $r_Q \ll \frac{1}{3} R_V$. Notice, that in the opposite to the absolute production rate (see Eq. (13)), in the regime of CT $\Sigma_\nu$ is insensitive to the wave function of the vector meson. Thus, predictions of attenuation effects are much less model dependent, and in this paper we concentrate on the nuclear transparency. (At a small and moderate $Q^2$, when $r_{FSI} \sim R_V$, the observable $\Sigma_\nu'$ for the production of the radially excited vector mesons $V'$ is extremely sensitive to the nodal structure of the wave function of the $V'$, which may lead $\Sigma_\nu' < 0$ and to the antishadowing phenomenon [1,5,6].) We predict that $1 - T_A$ scales with $r_Q^2$ [1], i.e., with $(Q^2 + m_V^2)$:

$$1 - T_A \propto \frac{A}{R_{ch}(A)^2 r_{FSI}^2} \sim \frac{A}{R_{ch}(A)^2} \frac{1}{Q^2 + m_V^2},$$

(16)

where $R_{ch}(A)$ is the charge radius of the nucleus. The simple law (16) holds at $1 - T_A \ll 1$. 


Before the comparison of the free-nucleon and the nuclear production, we emphasize that the mechanism of the free-nucleon reaction changes drastically from \( Q^2 \lesssim m_V^2 \to Q^2 \gg m_V^2 \). As Sakurai, Fraas and Shildknecht [16] have emphasized long ago, from the essentially kinematical considerations one finds a dominance of the longitudinal cross section at large \( Q^2 \),

\[
\frac{\sigma_L}{\sigma_T} \approx \frac{Q^2}{m_V^2},
\]

which agrees with the recent data from the E665 [4] and NMC [17] collaborations. Here we just mention that the relationship (17) holds for the nonrelativistic quarkonium, and the relativistic effects slow down the rise (17) [18].

The relativistic effects also slow down the rapid decrease (13) of \( m_T \) with \( Q^2 \) [18]. Specifically, because of the CT property of \( \sigma(r) \) large values of \( r \) are favored in the integrand of the production amplitude. The photon wavefunction \( \Psi_{\gamma^*}(r, z) \) admits large \( r \) if \( \varepsilon \) Eq. (7) is small, i.e., if either \( z \) or \( 1-z \) is small. Such an asymmetric quark pairs in the vector meson have a large intrinsic longitudinal momentum, which suppresses the wavefunction of the vector meson. Nonetheless, at very large \( Q^2 \) the production amplitude will be dominated by the asymmetric quark configurations in the vector meson, rather than by the nonrelativistic configurations \( z \sim 1/2 \). Even with allowance for the relativistic corrections, FSI effects in \( \sigma_L \) and \( \sigma_T \) differ little. We predict the approximately \( A \)-independent polarization density matrix of the produced \( \rho^0 \) mesons. The E665 data [4] confirm this. The detailed discussion of the relativistic effects in the \( \rho^0 \) production goes beyond the scope of the present paper and will be presented elsewhere [18]. Here we just emphasize that because \( \Psi_{\gamma^*}(r, z) \) is well understood [7], the virtual photoproduction offers a unique opportunity of scanning the wave function of vector mesons [1].

Regarding the experimental definition of the exclusive production, the word of caution is in order. The coherent production selects the truly exclusive production and its unambiguous signature is a very narrow diffraction peak. In the opposite to that, because of the limited energy resolution and imperfect rejection of the particle production in the target vertex, the incoherent sample is contaminated by inelastic interactions \( \gamma^*p \to Vp^* \) with excitation of the target proton. (Such an inelastic background was a major problem in earlier measurements of the exclusive \( \rho^0 \) production in deep inelastic scattering [19].) There are good reasons to expect that the ratio of the inelastic background to the exclusive reaction on the free and bound nucleons will be approximately the same. Also, in the regime of CT the cross section \( \sigma(r) \) has a universal dependence on \( r \), which is insensitive to the transition in the target nucleon vertex (for instance, see the considerations in [7,8]). Consequently, the inelastic background should only weakly affect the \( A \)- and \( Q^2 \)-dependence of nuclear transparency for the incoherent production. A comparison of coherent cross sections on two nuclei is ambiguity free. However, the relative normalization of the coherent nuclear cross section to the hydrogen cross section depends on the inelastic background in the latter. For this reason we compare \( \sigma_{coh}(A) \) and \( \sigma_{coh}(C) \).

In Fig.1 we show our predictions for the nuclear transparency for the incoherent cross section as a function of \( Q^2 \). (Predictions for production on the iron nucleus were published [1] before the E665 data were reported.) Nuclear attenuation is very strong at small \( Q^2 \) and gradually decreases with \( Q^2 \). This rise of nuclear transparency \( T_A \) with \( Q^2 \) is particularly dramatic for the heavy nuclei (\( Ca, Pb \)), and leaves no doubts that the onset of CT is observed.
Notice, that even at the highest $Q^2 \sim 10$ GeV$^2$ of the E665 experiment we predict rather strong attenuation effect: although at this value of $Q^2$ Eq. (1) gives $r_Q \approx 0.13$ f, the Eq. (13) shows that FSI is dominated by quark configuration having the relatively large transverse size $r_{FSI} \sim 0.5$ f. The large value of $r_{FSI}$ also shows that the relativistic effects are not yet important in the nuclear attenuation calculations. Our predictions for the $Q^2$ dependence of the nuclear transparency are in good agreement with the E665 data [4]. These data were taken at $\nu \approx 200$ GeV, so that the frozen size condition $l_e, l_f > R_A$ is fulfilled over the whole kinematical range of the E665 experiment.

In Fig. 2 we present our predictions for the $Q^2$ dependence of nuclear transparency for the forward coherent production on nuclei

$$T_A^{(coh)} = \left. \frac{d\sigma_A^{(coh)}}{A^2d\sigma_N} \right|_{\vec{q}^2=0}. \tag{18}$$

We predict a rise of $T_A^{(coh)}$ with $Q^2$. Even at $Q^2 \sim 10$ GeV$^2$ we predict a substantial departure from $T_A^{(coh)} = 1$ expected for the complete CT. The experimental determination of the absolute value of the forward production cross section is difficult as it requires high $\vec{q}^2$-resolution. Because of the above mentioned problem of the inelastic background in the hydrogen data we present predictions for the $A/C$ ratio.

The total coherent production cross section is much less sensitive to the $\vec{q}^2$-resolution. In Fig. 3 we present our predictions for the $Q^2$ dependence of the coherent production cross section relative to the cross section for the carbon nucleus. For the regime of complete CT and/or vanishing FSI,

$$R_{coh}^{(CT)}(A/C) = \frac{12\sigma_A}{A\sigma_C} \approx \frac{AR_{coh}(C)^2}{12R_{coh}(A)^2}, \tag{19}$$

which gives $R_{coh}^{(CT)}(Ca/C) = 1.56$ and $R_{coh}^{(CT)}(Pb/C) = 3.25$. We find a good quantitative agreement with the $Q^2$ dependence observed in the E665 experiment. As we have emphasized above, the coherent production is free of the inelastic background, and the E665 data on the Pb/C ratio give a particularly unambiguous evidence for the onset of CT.

The (approximate) $A^\alpha$ parametrization is a convenient short-hand representation of the $A$-dependence of nuclear cross sections. The so defined exponent $\alpha$ slightly depends on the range of the mass number $A$ used in the fit. Then, Eq. (13) predicts that $\alpha_{inc}(Q^2)$ and $\alpha_{coh}(\vec{q}^2 = 0)$ tend to 1 and 2 from below, as $Q^2$ increases. In the limit of vanishing final state interaction Eqs. (15,16,18) predicts $\sigma_{coh} \propto A^2/R_{ch}(A)^2 \sim A^{4/3}$, so that $\alpha_{coh}(Q^2)$ tends to $\approx \frac{4}{3}$ from below as $Q^2$ increases. In Fig. 4 we compare our estimate for the exponent $\alpha$ with the results of the E665 fits. (Our exponent $\alpha$ is defined as an average of the two values found from the ratio of theoretical prediction for the Pb/C and Ca/C cross section ratios; the uncertainties in $\alpha_{inc}$, $\alpha_{coh}(\vec{q}^2 = 0)$ and $\alpha_{coh}$ can be estimated as $\pm 0.03$, $\pm 0.03$ and $\pm 0.05$, respectively. The $A$-dependence of the no-FSI coherent cross section in the $C-Pb$ range of nuclei corresponds to the exponent $\alpha_{coh} \approx 1.39$ at $Q^2 \rightarrow \infty$.) Both the $\alpha_{coh}(Q^2)$ and $\alpha_{inc}(Q^2)$ rise with $Q^2$, which is still another way of stating that the E665 data confirm the onset of CT.

The scaling law (16) predicts that at $Q^2 \approx 9$ GeV$^2$ the nuclear attenuation for the incoherent $\rho^0$ production must be the same as for the incoherent real ($Q^2 = 0$) photoproduction
of the $J/\Psi$. The available $\rho^0$ and $J/\Psi$ production data were taken with somewhat different nuclear targets. For the $\rho^0$ production at $Q^2 = 7$ GeV$^2$ the E665 experiment gives $[T_{Pb}/T_C]_{\rho^0} = 0.6\pm0.25$. This can be compared with the NMC result $[T_{Sn}/T_C]_{J/\Psi} = 0.7\pm0.1$ for the real photoproduction of the $J/\Psi$ in the similar energy range. The substantial departure of the above values of $T_{A}/T_C$ from unity confirms our result [13,15] that even at $Q^2 \sim 10$ GeV$^2$ the FSI is controlled by a large $r_{FSI} \sim 0.5$ f and is not yet vanishing.

Similar scaling relationship holds for the coherent production of the $\rho^0$ and the $J/\Psi$. In the regime of complete CT Eqs. (11,12) give [6] $R_{coh}(Sn/C) = 2.76$, $R_{coh}(Fe/Be) = 2.82$, $R_{coh}(Pb/Be) = 4.79$. The experimental data on the real photoproduction of $J/\Psi$ give a solid evidence for nonvanishing FSI: $R_{coh}(Sn/C) = 2.15 \pm 0.10$ in the NMC experiment [21] and $R_{coh}(Fe/Be) = 2.28 \pm 0.32$, $R_{coh}(Pb/Be) = 3.47 \pm 0.50$ in the Fermilab E691 experiment [22]. In all cases the $\approx 25\%$ departure of the observed ratios for the $J/\Psi$ from predictions for the complete CT is of the same magnitude as in the highest $Q^2$ bin of the E665 data on the $\rho^0$ production (Fig.2). The analysis [5,6] within the same formalism as used in the present paper, gave a good quantitative description of the above $J/\Psi$ production data. Higher precision data on the $J/\Psi$ and $\rho^0$ production at higher $Q^2$ would be very interesting for further tests of our scaling law (14).

Our prediction [13] of a rapid decrease with $Q^2$ of the total production cross section which also follows from CT, is consistent with the experiment [17,19]. At very at large $Q^2$ the production cross section and the ratio $\sigma_L/\sigma_T$ do gradually become sensitive to the relativistic components of the wave function of vector mesons, but in view of Eq. (12) we predict a slow onset of the relativistic effects. As a matter of fact, the cross section (11) is well understood theoretically (for the discussion of how the small-$r$ behavior of $\sigma(r)$ is probed in deep inelastic scattering at small $x$ see [7,8]). Henceforth, one rather must use Eq. (5) for the $Q^2$-controlled scanning [1] and measurement of the wave function of vector mesons and for testing our understanding of these wave functions in the relativistic domain. To this end, the E665 results can be looked at as an important probe of $\sigma(r)$ from $r \sim R_V$ down to $r = r_{FSI} \sim 0.5$ f at $Q^2 = 7$ GeV$^2$. The observed attenuation corresponds to $\sigma(r = 0.5f) \sim 8$ mb, and the rise of the nuclear attenuation towards small $Q^2$ and $r \sim R_V$ proves that $\sigma(r)$ starts saturating at $\sigma(r) > \sigma_{tot}(\pi N)$ only at $r \gtrsim 1$ f, in agreement with the scenario [7].

We conclude that the solid signal of color transparency is seen in the exclusive $\rho^0$ production in deep inelastic scattering [4]. This is the first quantitative confirmation of color transparency ideas [9-11] in the hard scattering process (for the early evidence for CT in the $\pi^- A \rightarrow \pi^0 A^*$ charge exchange reaction see [23]). The E665 data are in good agreement with predictions of the lightcone approach to the virtual photoproduction of vector mesons developed by us [1,2,5,6]. Quite nontrivial feature of our approach is that the very mechanism of CT enforces rather large transverse size $r \sim r_{FSI}$ of the quark configurations which participate the intranuclear FSI when travelling through the nucleus. The scaling law (10) enables one to relate CT effects in the $\rho^0$ and the $J/\Psi$ production. The both $\rho^0$ and $J/\Psi$ sets of the experimental data show that even at $Q^2 \sim 10$ GeV$^2$ the residual FSI is not yet vanishing, and confirm our prediction of the slow onset of CT.

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Figure captions:

Fig. 1 - Predictions of nuclear transparency $T_A = \sigma_A/A\sigma_N$ for the incoherent exclusive production of $\rho^0$ mesons vs. the E665 data [4].

Fig. 2 - Predictions of nuclear transparency $T_A^{(coh)} / T_C^{(coh)} = [144d\sigma_A/A^2d\sigma_C]|_{q^2=0}$ for the forward coherent production of the $\rho^0$ mesons.

Fig. 3 - Predictions of the $Q^2$ dependence of the ratio of cross sections $R^{coh}(A/C) = 12\sigma_A/A\sigma_C$ for coherent production of the $\rho^0$ mesons vs. the E665 data [4]. The arrows indicate predictions for the complete CT at $Q^2 \to \infty$.

Fig. 4 - Predictions of the $Q^2$ dependence of exponents of parametrizations $\sigma_A^{(inc)} \propto A^{\alpha_{inc}}$, $[d\sigma_A^{(coh)}/dq^2]|_{q^2=0} \propto A^{\alpha_{coh}(q^2=0)}$ and $\sigma_A^{(coh)} \propto A^{\alpha_{coh}}$ vs. the E665 data [4].
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