IMPACT OF SATURATION ON S-CHANNEL HELICITY NONCONSERVATION FOR DIFRACTIVE VECTOR MESONS

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Abstract

As Glauber has shown in 1959, the spin-flip phenomena caused by the conventional spin-orbit interaction do vanish in the scattering off heavy, strongly absorbing nuclei. On the other hand, the origin of the s-channel helicity nonconservation (SCHNC) in diffractive DIS is of the origin different from simple spin-orbit interaction, and here we demonstrate that SCHNC in vector meson production survives strong absorption effects in nuclear targets. The intranuclear absorption often discussed in terms of the saturation effects introduces a new large scale $Q_A^2$ into the calculation of diffractive vector meson production amplitudes. Based on the color dipole approach, we show how the impact of the saturation scale $Q_A^2$ changes from the coherent to incoherent/quasifree diffractive vector mesons.

1 Introduction: the fate of spin-orbit interaction for heavy nuclei and the mechanism of SCHNC

Heavy nuclei are strongly absorbing targets. Whether the spin-flip effects in high energy scattering are washed out by this absorption or not is not an obvious issue which we address in this communication on an example of diffractive vector mesons.

The standard argument for vanishing of the spin-flip in elastic scattering of spin $\frac{1}{2}$ particles off strongly absorbing nuclei goes as follows: In the presence of the spin-orbit interaction the scattering amplitude $f = A + 2B\hat{s}n$, where $\hat{s}$ is the spin operator, $n$ is the normal to the scattering plane and the partial wave expansion of the helicity-non-flip, $f_0$, and the helicity-flip, $f_1$, amplitudes reads [1]

$$f_0 = \frac{1}{2ip}\sum_l \{(l + 1)[\exp(2i\delta^+_l) - 1] + l[\exp(2i\delta^-_l) - 1]\} P_l(\cos \theta),$$
$$f_1 = \frac{1}{2ip}\sum_l [\exp(2i\delta^+_l) - \exp(2i\delta^-_l)] P^+_l(\cos \theta), \quad (1)$$

where $\delta^\pm_j$ are the scattering phases for $j = l \pm \frac{1}{2}$. In the presence of strong absorption the scattering phases acquire large imaginary parts, $\exp(2i\delta^\pm_j) \to 0$, and, consequently, for the momentum transfer $\Delta$ within the diffraction cone $f_1/f_0 \to 0$. More detailed treatment for elastic scattering of protons off nuclei is found in Glauber’s lectures [2], the net result is that the small contribution to the spin-flip amplitudes comes only from the periphery of the nucleus, so that $f_1/f_0 \propto A^{-1/3}$, where $A$ is the mass number of a nucleus.
In high-energy QCD the diffractive production of vector mesons,

\[ \gamma^* p \rightarrow V p' \]

proceeds via the exchange of colorless system of gluons in the \( t \)-channel.

The fundamental property of such (multiple) gluon exchange is an exact conservation of the \( s \)-channel helicity of quarks, recall the similar property in the QED scattering of electrons in the Coulomb field [3]. None the less, QCD predicts a non-vanishing helicity flip in diffractive production of vector mesons off unpolarized nucleons [4, 5]: the origin of this SCHNC is in the subtle possibility that the sum of helicities of the quark-antiquark pair in the diagrams of fig. (1) can be unequal to the helicity of photons and vector mesons. In the non-relativistic case the pure S-wave deuteron with spin up consists of the spin-up proton and neutron. However, in the relativistic case while the longitudinal virtual photon contains the \( \bar{q}q \) pair with \( \lambda_q + \lambda_{\bar{q}} = 0 \), the transverse photon with helicity \( \lambda_\gamma = \pm 1 \) besides the \( \bar{q}q \) state with \( \lambda_q + \lambda_{\bar{q}} = \lambda_\gamma = \pm 1 \), also contains the state with \( \lambda_q + \lambda_{\bar{q}} = 0 \), in which the helicity of the photon is carried by the orbital angular momentum in the \( \bar{q}q \) system. Furthermore, it is precisely the state \( \lambda_q + \lambda_{\bar{q}} = 0 \) which gives the dominant contribution to the absorption of transverse photons and the proton SF \( F_{2p}(x, Q^2) \) in the Bjorken limit.

From the point of view of the vector meson production, it is important that the transverse and longitudinal \( \gamma^* \) and \( V \) share the intermediate \( \bar{q}q \) state with \( \lambda_q + \lambda_{\bar{q}} = 0 \), which allows the \( s \)-channel helicity non-conserving (SCHNC) transitions between the transverse (longitudinal) \( \gamma^* \) and longitudinal(transverse) vector meson \( V \). As a matter of fact, this mechanism of SCHNC does not require an applicability of pQCD, B.G. Zakharov was the first to introduce it in application to nucleon-nucleon scattering [6]. The theoretical predicition of energy-independent SCHNC in diffractive vector meson production has been confirmed experimentally at HERA [7, 8], the detailed comparison of the theory and experiment is found in [9].

Now notice that the argument about exact SCHC of quarks and antiquarks applies equally to a one-pomeron exchange in the scattering off a free nucleon and to multiple pomeron exchange in the scattering off a nuclear target. Then for a sufficiently high energy such that the lifetime of the \( \bar{q}q \) fluctuation of the photon, often referred to as the coherence time, and the formation time of the vector meson are larger than the radius.
of the nucleus $R_A$ (for instance, see [10, 11]), the above described origin of SCHNC must be equally at work for the nuclear and free-nucleon targets. In this communication we expand on this point, from the practical point of view one speaks of the values of

$$x = \frac{Q^2 + m_V^2}{2\nu} < x_A \approx 0.1 \cdot A^{-1/3}.$$  \hfill (2)

Another property of interactions with nuclei is the so-called saturation scale $Q_A$ which for partons with $x \approx x_A$ defines the transverse momentum below which their density is lowered by parton fusion effects [12, 13, 14]. It is interesting to see how the emergence of the saturation scale affects the $Q^2$-dependence of diffractive vector meson production, in particular, its SCHNC properties. Here one must compare the saturation scale $Q_A^2$ to the usual hard scale for diffractive vector meson production [15, 16]

$$\overline{Q}^2 \approx \frac{1}{4}(Q^2 + m_V^2).$$  \hfill (3)

The further presentation is organized as follows. In section 2 we start with the formulation of the color-dipole approach to calculation of helicity amplitudes for diffractive vector meson production. In section 3 we introduce the scanning radius and comment on the potential sensitivity to the short-distance behaviour of the wave function of vector mesons. In section 4 we start a discussion of nuclear effects on an example of coherent diffraction $\gamma^* A \rightarrow VA$ when the recoil nucleus remains in the ground state. We demonstrate how the dependence $\propto Q^{-4}$ which is common to all helicity amplitudes changes to $\propto \overline{Q}^2 Q_A^{-2}$ for $\overline{Q}^2 \lesssim Q_A^2$. Even stronger impact of saturation is found for incoherent diffractive vector mesons considered in section 5. In the Conclusions we summarize our principal findings and comment on the possibilities of the COMPASS experiment at CERN.

2 The free nucleon target

For the purposes of our discussion it is convenient to resort to the color dipole formalism: the production process depicted in fig. (1) factorizes into splitting of the photon into $q\bar{q}$ dipole way upstream the target, s-channel helicity conserving elastic scattering of the dipole off a target, and projection of the $q\bar{q}$ dipole onto the vector meson state. We restrict ourselves to the contribution from the $q\bar{q}$ Fock states of the vector meson which is a good approximation for $x \sim x_A$. The momentum-space calculation of the helicity amplitudes has been worked out time ago in [17, 9], the crucial ingredient in preserving the rotation invariance is the concept of the running polarization vector for the longitudinal vector mesons. Following [18], we make the Fourier transform to the color-dipole space and represent the helicity amplitudes $A_{fi}(x, \Delta)$, where $i = \lambda\gamma$ and $f = \lambda V$ are helicities of the initial state photon and the final state vector meson, respectively, in the color dipole factorization form

$$A_{fi}(x, \Delta) = \langle V_f\mid A_{q\bar{q}}(r, \Delta)\mid \gamma^*_i \rangle = i \int_0^1 dz \int d^2r \sigma(r, \Delta) \exp[i\frac{1}{2}(1 - 2z)(r\Delta)]I_{fi}(z, r),$$  \hfill (4)

where $\Delta$ is the transverse momentum transfer in the $\gamma^* \rightarrow V$ transition, $I_{fi}(z, r) = \Psi_{V,f}(z, r)\Psi_{\gamma^*,i}(z, r)$ and the summation over the helicities $\lambda, \overline{\lambda}$ of the intermediate $q\bar{q}$ pair.
is understood. The wave function $\Psi_{V,f}(z, r)$ of the final state vector meson contains the spin-orbital part and the "radial" wave functions, defined in terms of the vertex function $\Gamma_{V}(z, k)\bar{q}S_{\mu}qV_{\mu}$, where $S_{\mu}$ is the relevant Dirac structure and $V_{\mu}$ is the running polarization vector which must be so defined as to guarantee the rotational invariance for the fixed invariant mass $M$ of the lightcone $q\bar{q}$ Fock state of the vector meson [17, 9],

$$M^2 = \frac{m_f^2 + k^2}{z(1-z)},$$

where $k$ is the transverse momentum of the quark in the vector meson and $z$ and $(1-z)$ are fractions of the lightcone momentum of the vector meson carried by the quark and antiquark, respectively. In the momentum-space $\psi_{V}(z, k) \propto \Gamma_{V}(z, k)/(M^2 - m_f^2)$, the Fourier transform to the dipole space depends on the Dirac structure $S_{\mu}$ and the "radial" wave functions, defined in terms of the vertex function $\Gamma_{V}(z, k)\bar{q}S_{\mu}qV_{\mu}$. For the sake of simplicity, here we take $S_{\mu} = \gamma_{\mu}$, the exact form for the pure $S$ and $D$ wave states is found in [17, 9], the major conclusions on the impact of nuclear absorption do not depend on the exact form of $S_{\mu}$. If one defines the radial wave functions for the transverse ($T$) and longitudinal ($L$) vector mesons as

$$\psi_T(z, r) = \int d^2k \psi(z, k) \exp(ikr), \quad \psi_L(z, r) = \int d^2k M \psi(z, k) \exp(ikr),$$

then

$$I_{LL} = 4Qz^2(1-z)^2K_0(\varepsilon r)\psi_L(z, r)$$

$$I_{TT} = m_f^2K_0(\varepsilon r)\psi_T(z, r) - [z^2 + (1 - z)^2]\varepsilon K_1(\varepsilon r)\psi_T(z, r)$$

$$I_{LT} = -i2z(1-z)(1-2z)\psi_L(z, r)\varepsilon K_1(\varepsilon r)\frac{(er)}{r}$$

$$I_{TL} = -i2Qz(1-z)(1-2z)K_0(\varepsilon r)\psi_T(z, r)\frac{(V^*r)}{r}$$

$$I_{TT'} = 4z(1-z)\varepsilon K_1(\varepsilon r)\psi_T(z, r)\frac{(er)^2}{r^2}$$

where $\psi_T(z, r) = \partial\psi_T(z, r)/\partial r$, the polarization vectors $e$ and $V$ are for the transverse photon and vector meson, respectively, $\varepsilon^2 = z(1-z)Q^2 + m_f^2$, the Bessel functions $K_0(x)$ and $K_1 = K_1'(x)$ describe the lightcone wave function of the photon, $I_{TT}$ and $I_{TT'}$ describe the helicity-non-flip and double-helicity-flip production of transverse vector mesons by transverse photons, in the latter case $V^* = e$.

The off-forward generalization of the color-dipole scattering amplitude has been introduced in [19]

$$\sigma(r, \Delta) = \frac{2\pi}{3} \int \frac{d^2k}{(k - \frac{1}{2}\Delta)^2(k + \frac{1}{2}\Delta)^2} F(x, \kappa, \Delta) \alpha_s(\kappa^2) \times \left\{[1-\exp(i\kappa r)] \cdot [1-\exp(-i\kappa r)] - [1-\exp\left(\frac{1}{2}i\Delta r\right)] \cdot [1-\exp\left(-\frac{1}{2}i\Delta r\right)] \right\}$$

Here $F(x, \kappa, \Delta)$ is the off-forward unintegrated differential gluon structure function of the nucleon, the gross features of its $\Delta$-dependence are discussed in [19].

The analysis of [19] focused on the $LL$ and $TT$ amplitudes, in which case the dominant contribution comes from $z \sim \frac{1}{2}$ and corrections to the $\Delta$-dependence from the factor
\[ \exp\left[\frac{i}{2}(1-2z)(\Delta r)\right] \] in (4) can be neglected. This factor is crucial, though, for the helicity-flip transitions. For small dipoles and within the diffraction cone the leading components of the LT and TL amplitudes come from the second term in the expansion

\[ \exp\left(\frac{1}{2}i(1-2z)(\Delta r)\right) = 1 + \frac{1}{2}i(1-2z)(\Delta r) \] \hspace{1cm} (12)

so that upon the azimuthal averaging the effective integrands take the form

\[ I_{LT} = \frac{1}{2}z(1-z)(1-2z)^2\psi_L(z,r)\varepsilon r K_1(\varepsilon r)(e\Delta) \] \hspace{1cm} (13)

\[ I_{TL} = \frac{1}{2}Qz(1-z)(1-2z)^2K_0(\varepsilon r)r\psi'_T(z,r)(V^*\Delta) \] \hspace{1cm} (14)

In the case of the double-flip TT' amplitude \( e^2 = 0 \) and one needs to expand the integrand up to the terms \( \propto (r\Delta)^2 \):

\[ \exp\left(\frac{1}{2}i(1-2z)(\Delta r)\right) \left(\frac{er}{r^2}\sigma(x, r, \Delta)\right) \Rightarrow \frac{1}{60}(e\Delta)^2r^2[\sigma(\infty, 0) + (1-2z)^2\sigma(r, 0)] \] \hspace{1cm} (15)

so that the corresponding integrand of the double-flip amplitude will be of the form

\[ I_{TT'} = \frac{1}{15}z(1-z)\varepsilon r^2K_1(\varepsilon r)\psi'_T(z,r)(e\Delta)^2[\sigma(\infty, 0) + (1-2z)^2\sigma(r, 0)] \] \hspace{1cm} (16)

3 The scanning radius, hard scale expansion and sensitivity to the short distance wave function of the vector meson

Consider the pQCD regime of large \( Q^2 \). The useful representation for small color dipoles is [20]

\[ \sigma(r, 0) = \frac{\pi^2}{3}r^2\alpha_S(q^2)G(x, q^2), \quad q^2 \approx \frac{10}{r^2}. \] \hspace{1cm} (17)

Then, because of the exponential decrease, \( K_{0,1}(\varepsilon r) \propto \exp(-\varepsilon r) \), the amplitudes for the free nucleon target will be dominated the contribution from \( r = r_S \), where the scanning radius

\[ r_S \sim \frac{a_S}{\varepsilon} \approx \frac{a_S}{Q} = \frac{2a_S}{\sqrt{Q^2 + m_V^2}}, \quad a_S \approx 3. \] \hspace{1cm} (18)

The simplest case is that of the LL amplitude:

\[ A_{LL} \propto Qr_S^2\sigma(r_S, 0)K_0(a_S) \int dzz^2(1-z)^2\psi_L(z, r_S) \]

\[ \propto \frac{Q}{Q^2}\alpha_S(\overline{Q}^2)G(x, \overline{Q}^2) \propto \frac{QG(x, \overline{Q}^2)\alpha_S(\overline{Q}^2)}{(Q^2 + m_V^2)^2}. \] \hspace{1cm} (19)

The expansion in powers of the scanning radius \( r_S \) is an expansion in inverse powers of the hard scale \( \overline{Q} \). Here one factor \( 1/(Q^2 + m_V^2) \) is the same as in the Vector Dominance
Model, in the color dipole language it can be identified with the overlap of the photon and vector meson wave functions, the second factor \(1/(Q^2 + m_V^2)\) derives from the pQCD form (17) of the dipole cross section.

The helicity-flip amplitudes will be of the form

\[
\mathcal{A}_{LT} \propto (e\Delta) r_S^3 \sigma(r_S, 0) a_S K_1(a_S) \int dzz(1-z)(1-2z)^2 \psi_L(z, r_S)
\]

\[
\mathcal{A}_{TL} \propto (V^*\Delta) Q r_S^3 \sigma(r_S, 0) K_0(a_S) \int dzz(1-z)(1-2z)^2 \psi_T(z, r_S)
\]

Notice the sensitivity to the short-distance behavior of the vector meson wave function in the last result. The soft, oscillator-like interaction would give the wave function which at short distances is a smooth function of \(r^2\), to that

\[
\psi_T'(z, r) \sim -\frac{r}{R_V} \psi_T(z, 0)
\]

whereas the attractive Coulomb interaction at short distances suggests \([15, 21, 22]\) "hard", Coulomb-like \(\psi_T(z, r) \propto \exp(-r/R_C)\) when

\[
\psi_T'(z, r) \sim -\frac{1}{R_C} \psi_T(z, 0)
\]

In the case of the soft short-distance wave function the helicity-flip amplitude \(\mathcal{A}_{TL}\) would acquire extra small factor \(r_S/R_V \propto 1/(R_V Q)\). The similar discussion is relevant to the contribution to the \(\mathcal{A}_{TT}\) from the term \(-\varepsilon K_1(\varepsilon r)\psi_T'(z, r)\) in \(I_{TT}\), eq. (8), which for the hard short-distance wave function is larger and somewhat enhances the transverse cross section \(\sigma_T\) and lowers \(\sigma_L/\sigma_T\), see also the discussion in [23]. Finally, the leading term of expansion in powers of \(r_S\) of the double-helicity-flip amplitude is of the form

\[
\mathcal{A}_{TT} \propto (e\Delta)^2 z(1-z)r_S^3 a_S K_1(a_S) \sigma(\infty, 0) \int dzz(1-z)\psi_T'(z, r_S).
\]

It is proportional to the dipole cross section for the nonperturbative large color dipole [5]

## 4 Nuclear saturation effects in coherent diffraction

In the coherent diffractive production of vector mesons the target nucleus remains in the ground state,

\[\gamma^* A \rightarrow VA\]

For heavy nuclei such that their radii are much larger than the dipole size and the diffraction slope in the dipole-nucleon scattering, only the forward dipole-nucleon scattering enters the calculation of the nuclear profile function. Compared to the conventional derivation of the Glauber formulas for the nuclear profile functions, there are little subtleties with the presence of the phase factor \(\exp[-\frac{1}{2}(1-2z)(r\Delta)]\) in the color dipole factorization formula (4), but a careful rederivation gives the nuclear diffractive amplitudes of a form

\[
\mathcal{A}_{fi} = 2i \int d^2b \exp(-i\Delta b) |V_f| \left\{ 1 - \exp[-\frac{1}{2}\sigma(r, 0) T(b)] \right\} |\gamma_i^*\rangle \exp[i\frac{j}{2}(1-2z)(r\Delta)].
\]
where \( T(b) = \int dz n_A(b, z) \) is the standard nuclear optical thickness at an impact parameter \( b \). A comparison with the free nucleon amplitude (4) shows that the nuclear profile function

\[
\Gamma(b, r) = 1 - \exp[-\frac{1}{2} \sigma(r, 0) T(b)]
\]

plays the rôle of the color dipole cross section per unit area in the impact parameter space.

First of all, we notice that the color dipole dependence of the overlap of wave functions of the photon and vector meson does not change from the free nucleon to the nuclear case. The same is true of the calculation of expectation values over the orientation of color dipoles, see eqs. (14 - 16). The \( r \) dependence of the integrands changes, though.

Ref. [13] gives the detailed discussion of the reinterpretation of the nuclear profile function in terms of the saturating nuclear gluon density [12, 14] and of the limitations of such an interpretation for observables more complex than the single particle spectra. For the purposes of our discussion it is sufficient to know that in terms of the so-called saturation scale

\[
Q_A^2 = \frac{4\pi^2}{3} \alpha_s(Q_A^2) G(Q_A^2) T(b)
\]

the nuclear attenuation factor in (26) can be represented as

\[
\exp[-\frac{1}{2} \sigma(r, 0) T(b)] \approx \exp[-\frac{1}{8} Q_A^2 r^2]
\]

Let the nucleus be very heavy such that the saturation scale \( Q_A^2 \) is very large (for the estimates of \( Q_A^2 \) for realistic nuclei see [13]). For color dipoles with \( r^2 > r_A^2 = 8/Q_A^2 \) the nucleus is opaque, i.e., we have a saturated \( \Gamma(b, r) \approx 1 \) independent on the dipole size. The new large scale \( Q_A^2 \) must be compared to \( Q^2 \) of eq. (3).

First, there is a trivial case of \( Q^2 \gg Q_A^2 \). In this case \( r_S^2 \ll r_A^2 \), so that \( 2\Gamma(b, r) = \sigma(r_s, 0) T(b) \), i.e., the impulse approximation is at work and the nuclear amplitude has precisely the same structure as the free nucleon one,

\[
A_f^{(A)}(\Delta) = A_f^{(N)}(\Delta) \int d^2 b \exp(-ib\Delta) T(b) = A_f^{(N)}(\Delta) \cdot A \cdot G_{em}(\Delta),
\]

where \( G_{em}(\Delta) \) is the charge form factor of a nucleus.

Much more interesting is the case of coherent diffractive DIS at \( Q^2 \ll Q_A^2 \), when the color dipoles in the photon have the dipole size \( r \sim 1/Q \gg r_A^2 \), i.e., DIS proceeds in the regime of opacity and saturated color dipole cross section per unit area in the impact parameter space. Apart from this difference all the arguments of section 3 on the scanning radius will be fully applicable, only the scale for the scanning radius will change to \( a_S \approx 1 \). However, the, the power expansion in \( r_S \) will change: instead of the common factor \( r_S^4 \) of section 3 for the free nucleon target one obtains \( r_S^2 r_A^2 \) for diffraction off nuclei in the saturation scale, i.e., the substitution

\[
\left. \frac{1}{Q^2} \right|_N \Rightarrow \left. \frac{1}{Q^2 Q_A^2} \right|_A
\]

in the common prefactor of all helicity amplitudes. Otherwise there is no nuclear mass number dependent suppression of the relative strength of the helicity-flip and non-flip amplitudes compared to the free nucleon case. Finally, the \( \Delta \) dependence of the effective
nuclear form factor will be close to the ∆ dependence of an amplitude of elastic scattering on a black disc,
\[ G_{em}(\Delta) \Rightarrow \frac{J_1(R_A \Delta)}{R_A \Delta} \]  

We emphasize that although the presence of those nuclear form factors limit the practical observation of coherent diffractive DIS to momentum transfers within the nuclear diffraction cone, \( \Delta^2 \lesssim R_A^2 \), and these small ∆ cause the kinematical suppression of the helicity-flip within the coherent cone, there is no nuclear suppression of helicity flip even on a black nucleus.

5 Incoherent/quasielastic diffraction

In the incoherent (quasielastic, quasifree) diffractive vector meson production
\[ \gamma^* A \rightarrow V A^* \]
one sums over all excitations and breakup of the target nucleus without production of secondary particles. The process looks like a production off a quasi-free nucleon of the target subject to certain intranuclear distortions of the propagating dipoles. The relevant multichannel formalism has been worked out in [24], the generalization to the color dipole formalism for \( z \approx \frac{1}{2} \) is found in [10, 11]. Here we notice that in the color dipole language, the calculation of the helicity amplitudes will be exactly the same as for the free nucleon target but with the extra attenuation factor of eq. (28) in all the integrands. i.e., the incoherent differential cross section equals
\[ \frac{d\sigma}{d\Delta^2}(\gamma_i^* A \rightarrow V_f A^*) = \int d^2b T(b) \frac{d\sigma_{qel}}{d\Delta^2}, \]  

where the amplitudes of the quasielastic production off a quasi-free nucleon are given by
\[ A_{fi}(x, \Delta) = i \int_0^1 dz \int d^2 r \sigma(r, \Delta) \exp[-\frac{1}{2} \sigma(r, 0) T(b)] \exp[i(1-2z)(r\Delta)] I_{fi}(z, r), \]

In the genuine hard regime of \( Q^2 \gg Q_A^2 \) the nuclear attenuation can be neglected and one recovers the free nucleon cross section times the number of nucleons. In the opposite regime of strong saturation, \( Q^2 \ll Q_A^2 \), the \( r \) dependence of the attenuation factor is stronger than that of the photon wave functions \( K_{0,1}(\varepsilon r) \). Then, repeating the derivation of the scanning radius in section 3, one will find
\[ r_S^2 \approx \frac{3}{2} r_A^2. \]

The functional dependence of helicity amplitudes on the scanning radius \( r_S \) will be the same as for the free nucleon target with one exception. Namely, the Bessel functions in the photon wave function shall enter with the argument
\[ a_S = \varepsilon r_S \approx \frac{\sqrt{3} Q}{Q_A} \ll 1. \]
In this limit

$$K_0(a_S) \approx \log \left( \frac{Q_A}{Q} \right)$$  \hspace{1cm} (36)

which shows that some of the amplitudes will have a logarithmic enhancement, whereas

$$K_1(a_S) \approx \frac{Q_A}{Q}$$  \hspace{1cm} (37)

is indicative of even stronger enhancement. However, the closer inspection of the helicity-flip amplitude $A_{LT}$ (14) shows that $K_1(a_S)$ enters as a product $a_S K_1(a_S)$ which is a smooth function at $a_S \ll 1$. The helicity-flip amplitude $A_{TL}$ of eq. (14) exhibits only the weak logarithmic enhancement (36). The case of the helicity-non-flip $A_{TT}$ and double-flip $A_{TT'}$ is a bit more subtle. Here one encounters

$$-\varepsilon K_1(\varepsilon r_S) \psi'_T(z, r_S) \sim - \frac{1}{r_S} \psi'_T(z, r_S)$$  \hspace{1cm} (38)

which for the soft short-distance wave function can be estimated as

$$\frac{1}{R^2_V} \psi_T(z, r_S)$$  \hspace{1cm} (39)

whereas for the hard Coulomb wave function one finds an enhancement

$$\frac{Q_A}{R^2_C} \psi_T(z, r_S)$$  \hspace{1cm} (40)

To summarize, strong nuclear absorption does not generate any special suppression of the helicity-flip amplitudes compared to the non-flip ones. Furthermore, the estimate (40) suggests even a possibility of an enhancement of the double-flip transitions depending on the hardenss of the short-distance wave function of the vector meson. Finally, this discussion shows that in the saturation regime the saturation scale $Q^2_A$ becomes the hard factorization scale for incoherent diffractive production. Namely, for $Q^2 \ll Q^2_A$ this amounts to the substitution

$$\frac{1}{Q^2} \bigg|_N \Rightarrow \frac{1}{Q^2_A} \bigg|_A$$  \hspace{1cm} (41)

in the common prefactor of all helicity amplitudes. Similarly, the diffraction slope for the vector meson production will be the same as that for the free nucleon target but taken for the hard scale $Q_A$.

Here we focused on the single incoherent scattering approximation. The higher order incoherent interactions can readily be treated following the technique of ref. [24], they wouldn’t change major conclusions on the interplay of the DIS hard scales $Q^2$ and the saturation scale $Q^2_A$.

6 Conclusions

Our principal finding is a lack of nuclear suppression of the helicity-flip phenomena in hard diffractive production off strongly absorbing nuclei, which is in striking contrast to the familiar strong nuclear attenuation of the spin-orbit interaction effects as predicted
by the Glauber theory. The QCD mechanism behind this finding is that absorption only affects the color dipole-nucleus scattering amplitude in which the $s$-channel helicity of the quark and antiquark is anyway conserved exactly. The helicity flip originates from the relativistic mismatch of the sum of helicities of the quark and antiquark and the helicity of the vector meson and photon. Within the color dipole approach we demonstrated how the expansion of helicity amplitudes in powers of the scanning radius and saturation radius changes from the free nucleon to coherent nuclear to incoherent (quasielastic) nuclear diffractive production.

It is proper to recall the early claims by Greenberg and Miller of the so-called vector color transparency - the vanishing spin-flip phenomena in hard processes on nuclear targets [25]. These authors considered the polarization of ejected protons in quasielastic scattering $A(e,e'p)$ off nuclei in the version of spin-orbit interaction model and overlooked Zakharov’s helicity mismatch mechanism which, as we demonstrated above, persists in hard scattering. We discussed diffractive DIS, but all the arguments are equally applicable to the hadron-nucleus scattering.

The coherent and incoherent diffractive vector meson production off nuclei can be studied experimentally in the COMPASS experiment at CERN [26]. Whereas we are confident in our predictions for perturbatively large saturation scale $Q_A^2$, the numerical estimates for the saturation scale give disappointingly moderate $Q_A^2 \sim 1 \text{ GeV}^2$. None the less, the qualitative pattern of predicted changes of the $Q^2$ dependence of diffractive vector meson production from the free nucleon to coherent nuclear and incoherent nuclear cases must persist even at moderately large $Q_A^2$.

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