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SATURATION AND s-CHANNEL HELICITY NONCONSERVATION IN DIFFRACTIVE DIS

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The s-channel helicity conservation (SCHC) in diffractive DIS is known to break down despite the exact conservation of the s-channel helicity of quarks in high energy QCD. Here we demonstrate that the s-channel helicity nonconservation (SCHNC) in diffractive vector meson production survives strong absorption effects in nuclear targets. This must be contrasted to the classic 1959 result of Glauber that the spin-flip phenomena caused by the conventional spin-orbit interaction do vanish in the scattering off heavy, strongly absorbing nuclei. 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.

Сохранение s-канальной спиральности (ССКС) в дифракционном ГНР нарушается несмотря на точное сохранение s-канальной спиральности кварков в КХД при высоких энергиях. Мы демонстрируем, что несохранение s-канальной спиральности (НССКС) в дифракционном образовании векторных мезонов выживает при учете сильного поглощения в ядерных мишенях, что контрастирует с классическим результатом Глаубера (1959 г.) об отсутствии спин-орбитального взаимодействия при рассеянии на тяжелых сильно поглощающих ядрах. Внутриядерное поглощение, часто обсуждаемое в терминах насыщения, вносит в вычисление мезонных амплитуд новый большой масштаб Q_A^2. В рамках приближения цветных диполей мы показываем, как влияние Q_A^2 меняется от когерентных к некогерентным/квазисвободным дифракционным векторным мезонам.

INTRODUCTION

The spin-orbit interaction is of fundamental importance for the electron structure of atoms. Still in their 1954 classic study of high-energy electron scattering Yennie, Ravenhall and Wilson noticed that «for a spherically symmetric potential the scattering in a given direction will be the same for both spin orientations» [1] which is a manifestation of the s-channel helicity conservation (SCHC) of electrons in QED (for this interpretation of the YRW finding see the v. 4 of the Landau–Lifshits textbook [2]). In the realm of hadronic scattering, the 70's–80's were dominated by the Gilman et al. idea of SCHC in high-energy diffraction scattering ([3], for the review of early tests of SCHC see [4–6]). An explicit assumption of vanishing pomeron-exchange contribution to certain helicity amplitudes at high energies is behind the celebrated Burkhardt–Cottingham (BC) sum rule for the spin structure function $g_2(x, Q^2)$ [7]. However, the high-energy spin-physics proved to be much more fertile. It is by now well understood that the SCHC for quarks in high energy QCD does not entail the SCHC for hadrons because the helicity of hadrons differs from the sum of helicities of its constituent quarks the difference being due to the orbital angular momentum of constituents [8]. The explicit realization of this mechanism of s-channel helicity nonconservation (SCHNC) in diffractive DIS into continuum [9] and vector mesons has been published in [10–12]. This mechanism, in conjunction with the same mechanism for SCHNC for the Pomeron-nucleon coupling [8] and multipomeron-exchange unitarity corrections, was shown to break the assumptions behind the BC sum rule [13]. Similarly, the mutipomeron exchanges in DIS off deuterons give rise [14] to the tensor polarization of sea quarks in the deuteron which does not vanish at small x and invalidates the Close–Kumano sum rule [15]. A review of many aspects of the single- and double-spin asymmetries in high energy DIS, Drell–Yan and hadronic processes is found in [16, 17] and references therein. Here we focus on a still another new feature of SCHNC in diffractive DIS off nuclear targets — we demonstrate that in striking contrast to the familiar spin-orbit interaction effects it persists in the regime of strong nuclear absorption.

The further presentation is organized as follows. In Sec. 1 we explain briefly why the effect of the spin-orbit interaction vanishes in elastic scattering on a strongly absorbing target. In Sec. 2 we formulate the color-dipole approach to the calculation of helicity amplitudes for diffractive vector meson production. In Sec. 3 we introduce the scanning radius, describe the Q^2 dependence of helicity amplitudes in terms of an expansion in powers of the scanning radius originating from the color dipole cross section and the overlap of the lightcone color dipole wave functions of the photon and vector meson. In Sec. 4 we start a discussion of nuclear effects on an example of coherent diffraction $\gamma^* A \to VA$ when the recoil nucleus remains in the ground state. The nuclear effects in the black nucleus or the saturation regime are described by the new hard scale — the saturation scale Q_A^2 — and we demonstrate how the dependence $\propto \overline{Q}^{-4} \sim (Q^2 + m_V^2)^{-2}$ which is common to all helicity amplitudes, changes to the mixed $\propto \overline{Q}^{-2} Q_A^{-2}$ for $\overline{Q}^2 \lesssim Q_A^2$. Here the factor Q_A^{-2} describes the blackness of the target nucleus, whereas the factor $\overline{Q}^2$ still comes from the photon-vector meson overlap. Even stronger impact of saturation is found for quasielastic (incoherent) diffractive vector mesons considered in Sec. 5 — the saturation scale Q_A^2 becomes the sole hard scale for $\overline{Q}^2 \lesssim Q_A^2$. In the Conclusions we summarize our principal findings and comment on the possibilities of the COMPASS experiment at CERN.

This contribution is a humble tribute to Anatoly Vasil'evich Efremov, one of the world experts in high-energy spin phenomena, the author of such celebrated discoveries as the fundamental role of the axial anomaly in the helicity structure function $g_1(x, Q^2)$ of the nucleon [18] and handedness [19] and of many other important works on spin phenomena in hard reactions [20, 21], on the occasion of his 70's birthday.

1. THE FATE OF SPIN-ORBIT INTERACTION FOR HEAVY NUCLEI

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 here on an example of diffractive vector mesons.

The standard argument for vanishing of the spin-flip in elastic scattering of spin $1/2$ particles off strongly absorbing nuclei goes as follows: In the presence of the spin-orbit interaction the scattering amplitude $f = f_0 + 2f_1\hat{\mathbf{s}}\mathbf{n}$, where $\hat{\mathbf{s}}$ is the spin operator, $\mathbf{n}$ 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 [2]

$$\begin{aligned} f_0 &= \frac{1}{2ip}\sum_l \{(l+1)[\exp{(2i\delta_l^+)} - 1] + l[\exp{(2i\delta_l^-)} - 1]\} P_l(\cos\theta), \\ f_1 &= \frac{1}{2p}\sum_l [\exp{(2i\delta_l^+)} - \exp{(2i\delta_l^-)}] P_l^1(\cos\theta), \end{aligned} \tag{1}$$

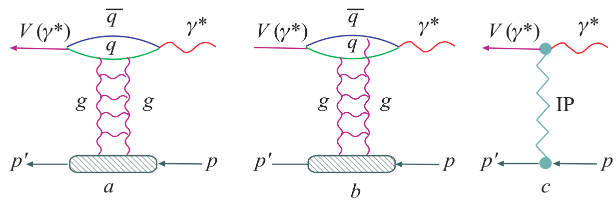
where $\delta_l^\pm$ are the scattering phases for $j = l \pm 1/2$. In the presence of strong absorption the scattering phases acquire large imaginary parts, $\exp{(2i\delta_l^\pm)} \to 0$, and, consequently, for the momentum transfer $\boldsymbol{\Delta}$ within the diffraction cone $f_1/f_0 \to 0$. More detailed treatment for elastic scattering of protons off nuclei is found in Glauber's lectures [22], 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.

2. SCHNC IN QCD

In high-energy QCD the diffractive production of vector mesons,

$$\gamma^* p \to V p',$$

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



a, b) The subset of two-gluon tower pQCD diagrams for the Pomeron exchange contribution (*c*) to the Compton scattering (DIS) and diffractive vector meson production. Not shown are two more diagrams with $q \leftrightarrow \bar{q}$

The fundamental property of such (multiple) gluon exchange is an exact conservation of the s-channel helicity of high-energy quarks. Nonetheless, QCD predicts a nonvanishing helicity flip in diffractive production of vector mesons off unpolarized nucleons [10, 11]: the origin of this SCHNC is in the subtle possibility that the sum of helicities of the quark-antiquark pair in the diagrams of the figure can be unequal to the helicity of photons and vector mesons. In the nonrelativistic 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 $q\bar{q}$ pair with $\lambda_q + \lambda_{\bar{q}} = 0$, the transverse photon with helicity $\lambda_\gamma = \pm 1$ besides the $q\bar{q}$ 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 $q\bar{q}$ 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 γ^* and V share the intermediate $q\bar{q}$ state with $\lambda_q + \lambda_{\bar{q}} = 0$, which allows the s-channel helicity nonconserving (SCHNC) transitions between the transverse (longitudinal) γ^* and longitudinal (transverse) vector meson $\mathbf{V}$. 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 [8]. The theoretical prediction of energy-independent SCHNC in diffractive vector meson production has been confirmed experimentally at HERA [23, 24], the detailed comparison of the theory and experiment is found in [25].

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 $q\bar{q}$ 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 [26, 27]), 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

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

Another property of interactions with nuclei is the so-called saturation scale Q_A which for partons with $x \lesssim x_A$ defines the transverse momentum below which their density is lowered by parton fusion effects [28–30]. 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 [31, 32]

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

3. 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 figure 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 [12, 25], the crucial ingredient in preserving the rotation invariance is the concept of the running polarization vector for the longitudinal vector mesons. Following [33], we make the Fourier transform to the color-dipole space and represent the helicity amplitudes $\mathcal{A}_{fi}(x, \mathbf{\Delta})$, 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

$$\mathcal{A}_{fi}(x, \mathbf{\Delta}) = \langle V_f | \mathcal{A}_{q\bar{q}}(\mathbf{r}, \mathbf{\Delta}) | \gamma_i^* \rangle = \\ = i \int_0^1 dz \int d^2\mathbf{r} \sigma(\mathbf{r}, \mathbf{\Delta}) \exp\left[\frac{i}{2}(1-2z)(\mathbf{r}\mathbf{\Delta})\right] I_{fi}(z, \mathbf{r}), \quad (4)$$

where $\mathbf{\Delta}$ is the transverse momentum transfer in the $\gamma^* \to V$ transition; $I_{fi}(z, \mathbf{r}) = \Psi_{V,f}^*(z, \mathbf{r}) \Psi_{\gamma^*,i}(z, \mathbf{r})$ and the summation over the helicities $\lambda, \overline{\lambda}$ of the intermediate $q\bar{q}$ pair is understood. The wave function $\Psi_{V,f}(z, \mathbf{r})$ of the final state vector meson contains the spin-orbital part and the «radial» wave functions, defined in terms of the vertex function $\Gamma_V(z, \mathbf{k}) \bar{q} S_\mu q V_\mu$, where S_μ is the relevant

Dirac structure and V_μ 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 [12, 25],

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

where $\mathbf{k}$ 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 $\psi_V(z,\mathbf{k}) \propto \Gamma_V(z,\mathbf{k})/(M^2 - m_V^2)$, the Fourier transform to the dipole space depends on the Dirac structure S_μ. 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 [12,25], the major conclusions on the impact of nuclear absorption on helicity flip do not depend on the exact form of S_μ. If one defines the radial wave functions for the transverse (T) and longitudinal (L) vector mesons as

$$\begin{aligned} \psi_T(z,\mathbf{r}) &= \int d^2\mathbf{k}\psi(z,\mathbf{k})\exp(i\mathbf{kr}), \\ \psi_L(z,\mathbf{r}) &= \int d^2\mathbf{k}M\psi(z,\mathbf{k})\exp(i\mathbf{kr}), \end{aligned} \tag{5}$$

then

$$I_{LL} = 4Qz^2(1-z)^2K_0(\varepsilon r)\psi_L(z,r), \tag{6}$$

$$I_{TT} = m_f^2K_0(\varepsilon r)\psi_T(z,r) - [z^2 + (1-z)^2]\varepsilon K_1(\varepsilon r)\psi_T'(z,r), \tag{7}$$

$$I_{LT} = -i2z(1-z)(1-2z)\psi_L(z,r)\varepsilon K_1(\varepsilon r)\frac{(\mathbf{er})}{r}, \tag{8}$$

$$I_{TL} = -i2Qz(1-z)(1-2z)K_0(\varepsilon r)\psi_T'(z,r)\frac{(\mathbf{V}^*\mathbf{r})}{r}, \tag{9}$$

$$I_{TT'} = 4z(1-z)\varepsilon K_1(\varepsilon r)\psi_T'(z,r)\frac{(\mathbf{er})^2}{r^2}, \tag{10}$$

where $\psi_T'(z,r) = \partial\psi_T(z,r)/\partial r$, the polarization vectors $\mathbf{e}$ and $\mathbf{V}$ 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 $K_1 = K_0'(x)$ 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 $\mathbf{V}^* = \mathbf{e}$ and we used $(\mathbf{V}^*\mathbf{e}) = \mathbf{e}^2 = 0$.

The off-forward generalization of the color-dipole scattering amplitude has been introduced in [34]

$$\sigma(\mathbf{r},\mathbf{\Delta}) = \frac{2\pi}{3}\int \frac{d^2\boldsymbol{\kappa}}{\left(\boldsymbol{\kappa}-\frac{1}{2}\mathbf{\Delta}\right)^2\left(\boldsymbol{\kappa}+\frac{1}{2}\mathbf{\Delta}\right)^2}\mathcal{F}(x,\boldsymbol{\kappa},\mathbf{\Delta})\alpha_S(\boldsymbol{\kappa}^2)\times \\ \times\Bigg\{[1-\exp(i\boldsymbol{\kappa}\mathbf{r})][1-\exp(-i\boldsymbol{\kappa}\mathbf{r})]- \\ -\left[1-\exp\left(\frac{1}{2}i\mathbf{\Delta}\mathbf{r}\right)\right]\left[1-\exp\left(-\frac{1}{2}i\mathbf{\Delta}\mathbf{r}\right)\right]\Bigg\}. \quad (11)$$

Here $\mathcal{F}(x,\boldsymbol{\kappa},\mathbf{\Delta})$ is the off-forward unintegrated differential gluon structure function of the nucleon, the gross features of its $\mathbf{\Delta}$-dependence are discussed in [34].

The analysis of [34] focused on the LL and TT amplitudes, in which case the dominant contribution comes from $z \sim 1/2$ and corrections to the $\mathbf{\Delta}$-dependence from the factor $\exp[i/2(1-2z)(\mathbf{r}\mathbf{\Delta})]$ in (4) can be neglected. This factor which is due to the Fermi motion of quarks in the vector meson 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

$$\exp\left(\frac{1}{2}i(1-2z)(\mathbf{\Delta}\mathbf{r})\right) = 1+\frac{1}{2}i(1-2z)(\mathbf{\Delta}\mathbf{r}) \quad (12)$$

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

$$I_{LT} = \frac{1}{2}z(1-z)(1-2z)^2\psi_L(z,r)\varepsilon r K_1(\varepsilon r)(\mathbf{e}\mathbf{\Delta}), \quad (13)$$

$$I_{TL} = \frac{1}{2}Qz(1-z)(1-2z)^2K_0(\varepsilon r)r\psi'_T(z,r)(\mathbf{V}^*\mathbf{\Delta}). \quad (14)$$

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

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

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

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

4. THE SCANNING RADIUS AND HARD SCALE EXPANSION

Consider the pQCD regime of large Q^2. The useful representation for small color dipoles is [35]

$$\sigma(\mathbf{r},0) \quad = \quad \frac{\pi^2}{3} r^2 \alpha_S(q^2) G(x,q^2), \quad q^2 \approx \frac{10}{r^2}. \tag{17}$$

Then, because of the exponential decrease, $K_{0,1}(\varepsilon r) \propto \exp(-\varepsilon r)$, the amplitudes for the free nucleon target will dominate the contribution from $r = r_S$, where the scanning radius [36]

$$r_S \sim \frac{a_S}{\varepsilon} \approx \frac{a_S}{\overline{Q}} = \frac{2a_S}{\sqrt{Q^2 + m_V^2}}, \quad a_S \approx 3. \tag{18}$$

The simplest case is that of the LL amplitude:

$$\mathcal{A}_{LL} \propto Q r_S^2 \sigma(r_S, 0) K_0(a_S) \int_0^1 dz z^2 (1-z)^2 \psi_L(z, r_S) \propto \\ \propto \frac{Q}{\overline{Q}^4} \alpha_S(\overline{Q}^2) G(x, \overline{Q}^2) \propto \frac{Q G(x, \overline{Q}^2) \alpha_S(\overline{Q}^2)}{(Q^2 + m_V^2)^2}. \tag{19}$$

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

$$\mathcal{A}_{LT} \propto (\mathbf{e}\boldsymbol{\Delta}) r_S^2 \sigma(r_S, 0) a_S K_1(a_S) \int_0^1 dz z (1-z)(1-2z)^2 \psi_L(z, r_S), \tag{20}$$

$$\mathcal{A}_{TL} \propto (\mathbf{V}^* \boldsymbol{\Delta}) Q r_S^3 \sigma(r_S, 0) K_0(a_S) \int_0^1 dz z (1-z)(1-2z)^2 \psi'_T(z, r_S). \tag{21}$$

Finally, the leading term of expansion in powers of r_S of the double-helicity-flip amplitude is of the form

$$\mathcal{A}_{TT'} \propto (\mathbf{e}\boldsymbol{\Delta})^2 r_S^3 a_S K_1(a_S) \sigma(\infty, 0) \int_0^1 dz z (1-z) \psi'_T(z, r_S). \tag{22}$$

Notice that it is proportional to the dipole cross section for the nonperturbative large color dipole and as such is of a manifestly nonperturbative origin [11].

5. THE SENSITIVITY TO THE SHORT DISTANCE WAVE FUNCTION OF THE VECTOR MESON

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 $\mathbf{r}^2$, so that

$$\psi'_T(z,r) \sim -\frac{r}{R_V^2}\psi_T(z,0) \tag{23}$$

whereas the attractive Coulomb interaction at short distances suggests [31, 37, 38] «hard», Coulomb-like $\psi_T(z,r) \propto \exp(-r/R_C)$ when

$$\psi'_T(z,r) \sim -\frac{1}{R_C}\psi_T(z,0). \tag{24}$$

In the case of the soft short-distance wave function the helicity-flip amplitude $\mathcal{A}_{TL}$ would acquire extra small factor $r_S/R_V \propto 1/(R_V\overline{Q})$. The similar discussion is relevant to the contribution to the $\mathcal{A}_{TT}$ from the term $\propto -\varepsilon K_1(\varepsilon r)\psi'_T(z,r)$ in I_{TT}, Eq. (8), which for the hard short-distance wave function is larger and somewhat enhances the transverse cross section σ_T and lowers σ_L/σ_T, see also the discussion in [39].

6. NUCLEAR SATURATION EFFECTS: COHERENT DIFFRACTION

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

$$\gamma^* A \to VA.$$

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 $\exp[i/2(1-2z)(\mathbf{r}\boldsymbol{\Delta})]$ in the color dipole factorization formula (4), but a careful rederivation gives the nuclear diffractive amplitudes of the form

$$\mathcal{A}_{fi} = 2i\int d^2\mathbf{b}\,\exp(-i\boldsymbol{\Delta}\mathbf{b})\langle V_f|\Gamma(\mathbf{b},\mathbf{r})\exp\left[\frac{i}{2}(1-2z)(\mathbf{r}\boldsymbol{\Delta})\right]|\gamma_i^*\rangle, \tag{25}$$

where

$$\Gamma(\mathbf{b},\mathbf{r}) = 1-\exp\left[-\frac{1}{2}\sigma(r,0)T(\mathbf{b})\right] \tag{26}$$

is the nuclear profile function and $T(\mathbf{b}) = \int dz n_A(\mathbf{b}, z)$ is the standard nuclear optical thickness at an impact parameter $\mathbf{b}$.

A comparison with the free nucleon amplitude (4) shows that the nuclear profile function $\Gamma(\mathbf{b}, \mathbf{r})$ can be regarded as the color dipole cross section per unit area in the impact parameter space. The color dipole dependence of the overlap of wave functions of the photon and vector meson and the calculation of expectation values over the orientation of color dipoles, see Eqs. (14)–(16), does not change from the free nucleon to the nuclear case. The r dependence of the integrands changes, though.

Ref. 29 gives the detailed discussion of the reinterpretation of the nuclear profile function in terms of the saturating nuclear gluon density [28, 30] 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 nuclear saturation scale,

$$Q_A^2(\mathbf{b}) = \frac{4\pi^2}{3}\alpha_S(Q_A^2)G(Q_A^2)T(\mathbf{b}) \propto A^{1/3}\,, \tag{27}$$

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

$$\exp\left[-\frac{1}{2}\sigma(r,0)T(\mathbf{b})\right] = \exp\left[-\frac{\sigma(r,0)}{\sigma(r_A(\mathbf{b}),0)}\right] \approx \exp\left[-\frac{1}{8}Q_A^2(\mathbf{b})r^2\right], \tag{28}$$

where $\sigma(r_A(\mathbf{b}), 0) = 2/T(\mathbf{b})$ is an implicit definition of $r_A(\mathbf{b})$, a useful approximation is $r_A^2(\mathbf{b}) \approx 8/Q_A^2(\mathbf{b})$. We denote by Q_A^2 the value of $Q_A^2(\mathbf{b})$ averaged over centrality $\mathbf{b}$, which is appropriate for typical DIS events. For estimates of Q_A^2 for realistic nuclei see [29]. The new large scale Q_A^2 must be compared to $\overline{Q}^2$ of Eq. (3), or the scanning radius r_S must be compared to the saturation radius r_A.

First, there is a trivial case of $\overline{Q}^2 \gg Q_A^2$. In this case the scanning radius is very small, $r_S^2 \ll r_A^2$, so that

$$2\Gamma(\mathbf{b}, \mathbf{r}) = \sigma(\mathbf{r}, 0)T(\mathbf{b}), \tag{29}$$

i.e., the nuclear attenuation effects can be neglected and the impulse approximation is at work. Then the nuclear amplitude has precisely the same structure as the free nucleon one,

$$\mathcal{A}_{fi}^{(A)}(\mathbf{\Delta}) = \mathcal{A}_{fi}^{(N)}(\mathbf{\Delta})\int d^2\mathbf{b}\exp\left(-i\mathbf{b}\mathbf{\Delta}\right)T(\mathbf{b}) = \mathcal{A}_{fi}^{(N)}(\mathbf{\Delta})AG_{\rm em}(\mathbf{\Delta}), \tag{30}$$

apart from the overall factors A and $G_{\rm em}(\mathbf{\Delta})$ — 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 $r \sim 1/Q \gg r_A$. It is the regime of nuclear opacity and $\Gamma(\mathbf{b},\mathbf{r}) \approx 1$ independent of the dipole size. When cast in the form

$$2\Gamma(\mathbf{b},\mathbf{r}) = \sigma(r_A(\mathbf{b}),0)T(\mathbf{b}) \tag{31}$$

and compared to (29), that can be interpreted as a saturation of color dipole cross section per unit area in the impact parameter space. Then, repeating the considerations leading to (19), one finds the nuclear amplitude

$$\begin{aligned} \mathcal{A}_{fi}^{(A)}(\mathbf{\Delta}) &\propto AQr_S^2\sigma(r_A,0)K_0(a_S)\int_0^1 dzz^2(1-z)^2\psi_L(z,r_S)\frac{J_1(R_A\Delta)}{R_A\Delta} \propto \\ &\propto \frac{AQ}{Q_A^2\overline{Q}^2}\frac{J_1(R_A\Delta)}{R_A\Delta}\alpha_S(Q_A^2)G(x,Q_A^2). \end{aligned} \tag{32}$$

Here R_A is the nuclear radius. First, all the arguments of Sec. 2 for the dominance by the contribution from $r \sim r_S = a_S/\overline{Q}$ will be applicable but because of the change of the r dependence of the integrand the scale a_S will change to

$$a_S|_{\rm coh} \approx 1.$$

Second, for the same reason the common prefactor of all helicity amplitudes, r_S^4, of Sec. 2 for the free nucleon target will change to $r_S^2 r_A^2$ for diffraction off nuclei in the saturation scale, i.e.,

$$\left.\frac{1}{\overline{Q}^2}\frac{1}{\overline{Q}^2}\right|_N \Rightarrow \left.\frac{1}{\overline{Q}^2}\frac{1}{Q_A^2}\right|_A. \tag{33}$$

Third, the nuclear mass number dependence, $A/Q_A^2 \propto A^{2/3}$, corresponds to that for elastic scattering off a black disc. Fourth, the $\mathbf{\Delta}$-dependence given for the impulse approximation by the nuclear charge form factor $G_{\rm em}(\mathbf{\Delta})$ changes to the familiar black disc form $J_1(R_A\Delta)/R_A\Delta$.

We emphasize that although the presence of those nuclear form factors limits the practical observation of coherent diffractive DIS to the momentum transfers within the nuclear diffraction cone, $\mathbf{\Delta}^2 \lesssim R_A^{-2}$, and these small $\mathbf{\Delta}$ cause the kinematical suppression of the helicity-flip within the coherent cone, there is no A-dependent nuclear suppression of helicity flip even on a black nucleus. The finite, A-independent, renormalization of the relative magnitude of different helicity amplitudes is possible, though, because of the change of the r dependence of the integrands of helicity amplitudes and the resulting change of the scale a_S for the scanning radius from $a_S \approx 3$ for the free nucleon to the A-independent $a_S \approx 1$ for the nuclear target.

7. NUCLEAR SATURATION EFFECTS: INCOHERENT/QUASIELASTIC DIFFRACTION

In the incoherent (quasielastic, quasifree) diffractive vector meson production,

$$\gamma^* A \to V A^*,$$

one sums over all excitations and break-up of the target nucleus without production of secondary particles. The process looks like a production off a quasifree nucleon of the target subject to certain intranuclear distortions of the propagating dipoles. The relevant multichannel formalism has been worked out in [40], the generalization to the color dipole formalism for $z \approx 1/2$ is found in [26, 27]. 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

$$\frac{d\sigma(\gamma_i^* A \to V_f A^*)}{d\mathbf{\Delta}^2} = \int d^2\mathbf{b} T(\mathbf{b}) \frac{d\sigma_{\rm qel}}{d\mathbf{\Delta}^2}, \tag{34}$$

where $d\sigma_{\rm qel}/d\mathbf{\Delta}^2 = |\mathcal{A}^{(\rm qel)}|^2/16\pi$ and the helicity amplitudes of the quasielastic production off a quasifree nucleon are given by

$$\mathcal{A}_{fi}^{(\rm qel)}(x, \mathbf{\Delta}) = i \int_0^1 dz \int d^2\mathbf{r} \sigma(\mathbf{r}, \mathbf{\Delta}) \exp\left[-\frac{1}{2}\sigma(r, 0) T(\mathbf{b})\right] \times \\ \times \exp\left[\frac{i}{2}(1-2z)(\mathbf{r}\mathbf{\Delta})\right] I_{fi}(z, \mathbf{r}). \tag{35}$$

These results hold for the momentum transfer $\mathbf{\Delta}$ beyond the diffractive peak for coherent diffraction off nucleus but within the diffraction cone for the free nucleon reaction.

In the genuine hard regime of $Q^2 \gg Q_A^2$, i.e., for $r_S^2 \ll r_A^2$, the nuclear attenuation can be neglected and one recovers the free nucleon cross section times the number of nucleons A. In the opposite regime of strong saturation, $\overline{Q}^2 \ll Q_A^2$, 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 Sec. 2, one will find

$$r_S^2 \approx \frac{3}{2} r_A^2. \tag{36}$$

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

$$a_S = \varepsilon r_S \approx \frac{\sqrt{3}\,\overline{Q}}{Q_A} \ll 1. \tag{37}$$

In this limit

$$K_0(a_S) \approx \log\left(\frac{Q_A}{\overline{Q}}\right) \tag{38}$$

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

$$K_1(a_S) \approx \frac{Q_A}{\overline{Q}} \tag{39}$$

is indicative of even stronger enhancement. However, the closer inspection of the helicity-flip amplitude $\mathcal{A}_{LT}$ (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 $\mathcal{A}_{TL}$ of Eq. (14) exhibits only the weak logarithmic enhancement (38). The case of the helicity-non-flip $\mathcal{A}_{TT}$ and double-flip $\mathcal{A}_{TT'}$ is a bit more subtle. Here one encounters

$$-\varepsilon K_1(\varepsilon r_S)\psi_T'(z, r_S) \sim -\frac{1}{r_S}\psi_T'(z, r_S) \tag{40}$$

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

$$\frac{1}{R_V^2}\psi_T(z, r_S), \tag{41}$$

whereas for the hard Coulomb wave function one finds an enhancement factor

$$\frac{Q_A}{R_C}\psi_T(z, r_S). \tag{42}$$

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 (42) suggests even a possibility of an enhancement of the double-flip transitions depending on the hardness of the short-distance wave function of the vector meson. Finally, this discussion shows that in the saturation regime precisely the saturation scale Q_A^2 sets the hard factorization scale for incoherent diffractive production. Namely, for $\overline{Q^2} \ll Q_A^2$ this amounts to the substitution

$$\left.\frac{1}{\overline{Q}^4}\right|_N \Rightarrow \left.\frac{1}{Q_A^4}\right|_A \tag{43}$$

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. 40, they wouldn't change major conclusions on the interplay of the DIS hard scales $\overline{Q}^2$ and the saturation scale Q_A^2.

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 and to incoherent (quasielastic) nuclear diffractive production.

The coherent and incoherent diffractive vector meson production off nuclei can be studied experimentally in the COMPASS experiment at CERN [41]. 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 $Q_A^2 \sim 1$ GeV^2. Nonetheless, 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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