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DIFFRACTIVE VECTOR MESONS BEYOND THE S-CHANNEL HELICITY CONSERVATION

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We derive a full set, and determine the twist, of helicity amplitudes for diffractive production of light to heavy vector mesons in deep inelastic scattering. For large Q^2 all helicity amplitudes but the double-flip are calculable in perturbative QCD and are proportional to the gluon structure function of the proton at a similar hard scale. We find a substantial breaking of the s-channel helicity conservation which must persist also in real photoproduction.

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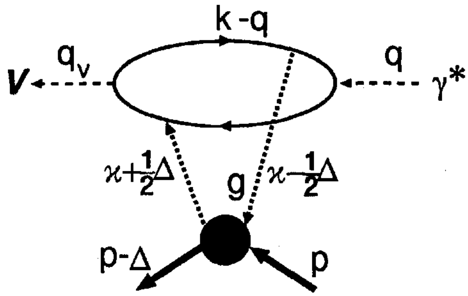
Diffractive virtual photoproduction of vector mesons $\gamma^* + p \to V + p'$, $V = \rho_0, \omega, \phi, J/\Psi, \Upsilon$, in deep inelastic scattering (DIS) at small $x = (Q^2 + m_V^2)/(W^2 + Q^2)$ is a testing ground of ideas on the QCD pomeron exchange and light-cone wave function (LCWF) of vector mesons ([1–5], for the recent review see [6]). (For the kinematics see Figure, $Q^2 = -q^2$ and $W^2 = (p+q)^2$ are standard DIS variables). It is a self-analyzing process because the helicity amplitudes can be inferred from vector meson decay angular distributions [7]. The property of s-channel helicity conservation (SCHC) and the dominance of transitions $\gamma_L^* \to V_L$ and $\gamma_T^* \to V_T$ with $R = \sigma_L/\sigma_T \approx Q^2/m_V^2$ are shared by nearly all models (L and T stand for the longitudinal and transverse polarizations). The nonperturbative contributions to σ_T tame the rise of R with Q^2 [2, 4], but still the theoretical predictions seem to exceed systematically the experimental evaluations of R. However, these experimental data analyses suffer from an unwarranted assumption of an exact SCHC [6], which needs a theoretical scrutiny. We report here a derivation of the full set of helicity amplitudes for transitions $\gamma_L^* \to V_L$, $\gamma_L^* \to V_T$, $\gamma_T^* \to V_T$, $\gamma_T^* \to V_L$ for all flavours and small to moderate momentum transfer $\mathbf{\Delta}$ within the diffraction cone. We find substantial s-channel helicity non-conserving (SCHNC) effects similar to the SCHNC LT interference found earlier by Pronyaev and two of the present authors (NNN and BGZ) for diffraction $\gamma^* p \to p' X$ into continuum states X [8].

The leading $\log\frac{1}{x}$ (LL$\frac{1}{x}$) pQCD diagrams for vector meson production are shown in Figure. We treat vector mesons as $q\bar{q}$ states with the $V\bar{q}q$ vertex $\Gamma_V V_\mu \bar{u}\gamma_\mu u$. At small x it is sufficient to compute the imaginary part of the amplitude, the real part is a small correction which can readily be reconstructed from analyticity [4] and we do not discuss it any more. We use the standard Sudakov expansion of all the momenta in the two lightcone vectors

$$p' = p - q\frac{p^2}{s}, \quad q' = q + p'\frac{Q^2}{s},$$



One of the four Feynman diagrams for the vector meson production $\gamma^* p \to V p'$ via QCD two-gluon pomeron exchange

such that $q'^2 = p'^2 = 0$ and $s = 2p' \cdot q'$, and the two-dimensional transverse component: $k = zq' + yp' + k_\perp$, $\kappa = \alpha q' + \beta p' + \kappa_\perp$, $\Delta = \gamma p' + \delta q' + \Delta_\perp$ and for the final vector meson $q_V = q + \Delta = q' + \frac{m_V^2 + \mathbf{\Delta}^2}{s} p' + \Delta_\perp$ (hereafter $\mathbf{k}, \mathbf{\Delta}, ..$ always stand for $k_\perp, \Delta_\perp$ etc.). As usual, only the so-called nonsense components in the Gribov's decomposition of gluon propagators do contribute in the high-energy limit,

$$D_{\mu\nu}(k) = \frac{2p'_\mu q'_\nu}{sk^2},$$

so that in the upper blob the amplitude $R_{\mu\nu\rho\sigma}$ of the subprocess $g_\mu \gamma^*_\sigma \to q\bar{q} \to g'_\nu V_\rho$ enters in the form $I(\gamma^* \to V) \propto p'_\mu p'_\nu R_{\mu\nu\rho\sigma} V^*_\rho e_\sigma$. The vertex function Γ is related to the radial LCWF of the $q\bar{q}$ Fock state of the vector meson as

$$\psi_V(z, \mathbf{k}) = \frac{\Gamma_V(z, \mathbf{k})}{D(m_V^2, z, \mathbf{k})}, \tag{1}$$

where $D(\mathbf{k}^2) = \mathbf{k}^2 + m_q^2 - z(1-z)m_V^2$. For the parameterizations of LCWF's of vector mesons based on the technique [9], see [4]. The corresponding quantity for photons, $\psi_\gamma(z, \mathbf{k}) = 1/D(-Q^2, z, \mathbf{k})$ only differs by the substitution $\Gamma_\gamma(z, \mathbf{k}) = 1$ and $m_V^2 \to -Q^2$. To the $LL\frac{1}{x}$ the lower blob is related to the unintegrated gluon density matrix $\mathcal{F}(x, \boldsymbol{\kappa}, \mathbf{\Delta})$ [5, 10, 11]. After standard elimination of the Sudakov parameters y, β etc. from the on-mass shell condition, for instance, $(k-q)^2 = -sy(1-z) - \mathbf{k}^2 - Q^2(1-z) = m_q^2$, the virtual photoproduction amplitude takes the form

$$A(x, Q^2, \mathbf{\Delta}) = is\frac{C_F N_c c_V \sqrt{4\pi\alpha_{em}}}{2\pi^2} \times$$

$$\times \int_0^1 dz \int d^2\mathbf{k} \int \frac{d^2\boldsymbol{\kappa}}{\kappa^4} \alpha_S(\max\{\kappa^2, \mathbf{k}^2 + m_V^2\}) \mathcal{F}(x, \boldsymbol{\kappa}, \mathbf{\Delta}) I(\gamma^* \to V), \tag{2}$$

where $N_c = 3$ is the number of colors, $C_F = (N_c^2 - 1)/2N_c$ is the Casimir operator, $C_V = 1/\sqrt{2}, 1/3\sqrt{2}, 1/3, 2/3$ for the $\rho^0, \omega^0, \phi^0, J/\Psi$ mesons, α_S and α_{em} are the strong and electromagnetic couplings, respectively. The dependence of $\mathcal{F}(x, \boldsymbol{\kappa}, \mathbf{\Delta})$ on the variable $\mathbf{\Delta}\boldsymbol{\kappa}$ corresponds to the subleading BFKL singularities [11] and can be neglected. For small $\mathbf{\Delta}$ within the diffraction cone

$$\mathcal{F}(x, \boldsymbol{\kappa}, \mathbf{\Delta}) = \frac{\partial G(x, \kappa^2)}{\partial \log \kappa^2} \exp(-\frac{1}{2} B_{3I\!P} \mathbf{\Delta}^2), \tag{3}$$

where $\partial G/\partial \log \kappa^2$ is the conventional unintegrated gluon structure function and, modulo to a slow Regge growth, the diffraction cone $B_{3\mathrm{IP}} \sim 6\ \mathrm{GeV}^{-2}$ [5]. The familiar polarization vectors for virtual photons are

$$e_L = \frac{1}{Q}\left(q' + \frac{Q^2}{s}p'\right), \quad e_T = e_\perp,$$

whereas for vector mesons

$$V_T = V_\perp + \frac{2(\mathbf{V}_\perp \cdot \mathbf{\Delta})}{s}(p' - q'), \quad V_L = \frac{1}{m_V}\left(q' + \frac{\mathbf{\Delta}^2 - m_V^2}{s}p' + \Delta_\perp\right), \tag{4}$$

such that $(V_T V_L) = (V_T q_V) = (V_L q_V) = 0$. Because of the small factor $1/s$ it is tempting to neglect the component $\propto p'$ in V_T but that would have been entirely erroneous. Indeed, closer inspection of the evaluation of $I(\gamma^* \to V)$ shows that $p'_\mu p'_\nu R_{\mu\nu\rho\sigma} p'_\rho e_\sigma \propto$ $\propto s p'_\mu p'_\nu R_{\mu\nu\rho\sigma} V^*_{\perp,\rho} e_\sigma$. The same is true of the contribution from the component $\propto \Delta_\perp$ in V_L.

It is convenient to define

$$\mathbf{\Psi}_{\gamma,V}(\mathbf{k}) = \mathbf{k}\psi_{\gamma,V}(z,\mathbf{k})\,, \chi_\gamma(z,\mathbf{k}) = \left[\mathbf{k}^2 + m_q^2 - z(1-z)Q^2\right]\psi_\gamma(z,\mathbf{k}),$$

$$\chi_V(z,\mathbf{k}) = \left[\mathbf{k}^2 + m_q^2 + z(1-z)m_V^2\right]\psi_V(z,\mathbf{k})$$

and

$$\eta_\gamma = \psi_\gamma(z,\mathbf{k}) - \psi_\gamma(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}\mathbf{\Delta})\,, \tag{5}$$

$$\eta_V = \psi_V(z,\mathbf{k} + (1-z)\mathbf{\Delta}) - \psi_V(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}(1-2z)\mathbf{\Delta})\,, \tag{6}$$

$$\varphi_\gamma = \chi_\gamma(z,\mathbf{k}) - \chi_\gamma(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}\mathbf{\Delta}) = -2z(1-z)Q^2\eta_\gamma\,, \tag{7}$$

$$\varphi_V = \chi_V(z,\mathbf{k} + (1-z)\mathbf{\Delta}) - \chi_V(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}(1-2z)\mathbf{\Delta}), \tag{8}$$

$$\boldsymbol{\phi}_\gamma = \mathbf{\Psi}_\gamma(z,\mathbf{k}) - \mathbf{\Psi}_\gamma(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}\mathbf{\Delta})\,, \tag{9}$$

$$\boldsymbol{\phi}_V = \mathbf{\Psi}_V(z,\mathbf{k} + (1-z)\mathbf{\Delta}) - \mathbf{\Psi}_V(z,\mathbf{k} - \boldsymbol{\kappa} + \frac{1}{2}(1-2z)\mathbf{\Delta})\,, \tag{10}$$

$$\begin{aligned}\Phi_2 &= \psi_\gamma(z,\mathbf{k} - (1-z)\mathbf{\Delta}) - \psi_\gamma(z,\mathbf{k} - \boldsymbol{\kappa} - \frac{1}{2}(1-2z)\mathbf{\Delta}), \\ &-\psi_\gamma(z,\mathbf{k} + \boldsymbol{\kappa} - \frac{1}{2}(1-2z)\mathbf{\Delta}) + \psi_\gamma(z,\mathbf{k} + z\mathbf{\Delta})\,,\end{aligned} \tag{11}$$

$$\begin{aligned}\mathbf{\Phi}_1 &= \mathbf{\Psi}_\gamma(z,\mathbf{k} - (1-z)\mathbf{\Delta}) - \mathbf{\Psi}_\gamma(z,\mathbf{k} - \boldsymbol{\kappa} - \frac{1}{2}(1-2z)\mathbf{\Delta}), \\ &-\mathbf{\Psi}_\gamma(z,\mathbf{k} + \boldsymbol{\kappa} - \frac{1}{2}(1-2z)\mathbf{\Delta}) + \mathbf{\Psi}_\gamma(z,\mathbf{k} + z\mathbf{\Delta})\,.\end{aligned} \tag{12}$$

Then the integrands $I(\gamma^* \to V)$ in (2) can be cast in the form which emphasizes the $V \leftrightarrow \gamma$ symmetry nicely:

$$I(\gamma_L^* \to V_L) = \frac{1}{Qm_V}\varphi_\gamma\varphi_V = -\frac{2Q}{m_V}z(1-z)\Phi_2\left[m_q^2 + \mathbf{k}^2 + z(1-z)m_V^2\right]\psi_V(z,\mathbf{k})\,, \tag{13}$$

$$I(\gamma_T^* \to V_T) = (\mathbf{V}^*\mathbf{e})[m_q^2\eta_V\eta_\gamma + (\boldsymbol{\phi}_V\boldsymbol{\phi}_\gamma)] +$$
$$+(1-2z)^2(\boldsymbol{\phi}_\gamma\mathbf{e})(\boldsymbol{\phi}_V\mathbf{V}^*) - (\boldsymbol{\phi}_V\mathbf{e})(\boldsymbol{\phi}_\gamma\mathbf{V}^*) =$$
$$= \left\{(\mathbf{V}^*\mathbf{e})[m_q^2\Phi_2 + (\mathbf{k}\boldsymbol{\Phi}_1)] + (1-2z)^2(\mathbf{V}^*\mathbf{k})(\mathbf{e}\boldsymbol{\Phi}_1) - (\mathbf{e}\mathbf{k})(\mathbf{V}^*\boldsymbol{\Phi}_1)\right\}\psi_V(z,\mathbf{k}), \qquad (14)$$

$$I(\gamma_T^* \to V_L) = \frac{(1-2z)}{m_V}(\boldsymbol{\phi}_\gamma\mathbf{e})\varphi_V =$$
$$= \frac{(1-2z)}{m_V}(\mathbf{e}\boldsymbol{\Phi}_1)\left[m_q^2 + \mathbf{k}^2 + z(1-z)m_V^2\right]\psi_V(z,\mathbf{k}), \qquad (15)$$

$$I(\gamma_L^* \to V_T) = \frac{(1-2z)}{Q}\varphi_\gamma(\mathbf{V}^*\boldsymbol{\phi}_V) = -2Qz(1-z)(1-2z)(\mathbf{V}^*\mathbf{k})\Phi_2\psi_V(z,\mathbf{k}). \qquad (16)$$

We consistently keep the quark mass so that our derivation holds from light to heavy vector mesons.

After simple shifts of the integration variables to the common argument of the vector meson LCWF, the $I(\gamma^* \to V)$ can also be cast in the second form shown in (13) – (16). This second form emphasizes the close analogy between amplitudes of vector meson production and diffraction into continuum discussed in [8, 10, 12], the major difference being the emergence of $\psi_V(z,\mathbf{k})$ instead of the plane WF. The second form is also convenient for the derivation of the $1/Q^2$ expansion, i.e., of the twist, of helicity amplitudes. To this end we notice that $\psi_V(z,\mathbf{k})$ varies on a hadronic scale $k^2 \sim R_V^{-2}$, where R_V is a radius of the vector meson, whereas $\psi_\gamma(z,\mathbf{k})$ is a slow function which varies on a large pQCD scale

$$\overline{Q}^2 = m_q^2 + z(1-z)Q^2, \qquad (17)$$

which allows a systematic twist expansion in powers of $1/\overline{Q}^2$ for Δ within the diffraction cone. In all cases but the double helicity flip the dominant twist amplitudes come from the leading $\log\overline{Q}^2$ ($LL\overline{Q}^2$) region of $\mathbf{k}^2 \sim R_V^{-2}, \boldsymbol{\Delta}^2 \ll \boldsymbol{\kappa}^2 \ll \overline{Q}^2$. After some algebra in (13) – (16) we find

$$I(\gamma_L^* \to V_L) = -\frac{Q}{m_V}\frac{4z(1-z)[m_q^2 + \mathbf{k}^2 + z(1-z)m_V^2]}{\overline{Q}^4}\cdot\psi_V(z,\mathbf{k})\boldsymbol{\kappa}^2, \qquad (18)$$

$$I(\gamma_T^* \to V_T; \lambda_V = \lambda_\gamma) = 2(\mathbf{V}^*\mathbf{e})\frac{m_q^2 + 2[z^2 + (1-z)^2]\mathbf{k}^2}{\overline{Q}^4}\psi_V(z,\mathbf{k})\boldsymbol{\kappa}^2, \qquad (19)$$

$$I(\gamma_L^* \to V_T) = -8\frac{(\mathbf{V}^*\boldsymbol{\Delta})Q}{\overline{Q}^2}\frac{z(1-z)(1-2z)^2}{\overline{Q}^4}\mathbf{k}^2\psi_V(z,\mathbf{k})\boldsymbol{\kappa}^2, \qquad (20)$$

$$I(\gamma_T^* \to V_L) = -2\frac{(\mathbf{e}\boldsymbol{\Delta})}{m_V}\frac{[m_q^2 + \mathbf{k}^2 + z(1-z)m_V^2](1-2z)^2}{\overline{Q}^4}\psi_V(z,\mathbf{k})\boldsymbol{\kappa}^2. \qquad (21)$$

Here we split explicitly the amplitude (14) into the SCHC helicity-non-flip component (19) and SCHNC double-helicity flip component

$$I(\gamma_T^* \to V_{T'}; \lambda_V = -\lambda_\gamma) = 2(\mathbf{V}^*\boldsymbol{\Delta})(\mathbf{e}\boldsymbol{\Delta})\frac{z(1-z)}{\overline{Q}^4}\left[6(1-2z)^2\frac{\boldsymbol{\kappa}^2}{\overline{Q}^2} + 1\right]\mathbf{k}^2\psi_V(z,\mathbf{k}). \qquad (22)$$

In all cases when $I(\gamma^* \to V) \propto \kappa^2$ the LL$\overline{Q}^2$ approximation is at work, the gluon structure function enters the integrand in the form

$$\int^{\overline{Q}^2} \frac{d\kappa^2}{\kappa^2} \frac{\partial G(x,\kappa^2)}{\partial \log \kappa^2} = G(x,\overline{Q}^2) \tag{23}$$

and after the z integration one finds that the helicity amplitudes will be proportional to $\alpha_S(Q_V^2)G(x,Q_V^2)$, where the pQCD hard scale

$$Q_V^2 \sim (0.1-0.2)(Q^2+m_V^2) \tag{24}$$

can depend on helicities because of the different end-point contributions from $z \to 0$ and $z \to 1$ where the running hardness $\overline{Q}^2$ is small [2, 4]. This issue and the sensitivity of helicity amplitudes to the LCWF of vector mesons will be discussed elsewhere. Here we cite our results for the dominant twist SCHC and single-flip SCHNC amplitudes $A_{\lambda_V \lambda_\gamma}$ in the helicity basis:

$$A_{0L} \propto \frac{Q}{m_V}\frac{\alpha_S(Q_V^2)}{(Q^2+m_V^2)^2}G(x,Q_V^2)\,, \tag{25}$$

$$A_{\pm\pm} \propto \frac{\alpha_S(Q_V^2)}{(Q^2+m_V^2)^2}G(x,Q_V^2)\,, \tag{26}$$

$$A_{\pm L} \propto \frac{\Delta}{m_V}\frac{\alpha_S(Q_V^2)}{(Q^2+m_V^2)^2}G(x,Q_V^2)\,, \tag{27}$$

$$A_{0\pm} \propto \frac{Q}{m_V}\frac{m_V\Delta}{Q^2+m_V^2}\frac{\alpha_S(Q_V^2)}{(Q^2+m_V^2)^2}G(x,Q_V^2)\,, \tag{28}$$

First, all above amplitudes are pQCD calculable for large Q^2 and/or heavy vector mesons. Second, the factor Q/m_V in the longitudinal photon amplitudes (25) and (27) is a generic consequence of gauge invariance irrespective of the detailed production dynamics. Third, the longitudinal Fermi momentum of quarks is $k_z \sim 1/2m_V(2z-1)$ and eqs. (20),(21) make it obvious that single-helicity flip requires a longitudinal Fermi motion of quarks and vanishes in the nonrelativistic limit. Similarly, double-helicity flip requires the transverse Fermi motion of quarks. Fourth, the leading SCHNC effect is an interference of A_{0L} and $A_{0\pm}$ and the first experimental indications for that have been reported recently [13]. We hotice here that a duality consistency [14] holds between the twist of $\propto A_{0L}A_{0\pm}$ and that of the LT interference structure function F_{LT}^D derived in [8]. The issue of duality for SCHNC diffractive DIS will be discussed elsewhere. Fifth, excitation of transverse mesons by longitudinal photons is of higher twist compared to excitation of longitudinal mesons by transverse photons.

In contrast to the above, the dominant twist double-helicity flip amplitude comes from the non-leading-log$\overline{Q}^2$ term 1 in the square brackets in (22) and is proportional to

$$\int^{\overline{Q}^2} \frac{d\kappa^2}{\kappa^4} \frac{\partial G(x,\kappa^2)}{\partial \log \kappa^2} \sim \frac{1}{\mu_G^2}G(x,\mu_G^2) \tag{29}$$

where a soft scale $\mu_G \sim$ 0.7-1 GeV is set by the inverse radius of propagation of perturbative gluons. Precisely the same nonperturbative quantity (29) describes the contribution from the $\gamma^* \to V$ transition vertex to the diffraction slope for helicity-non-flip amplitudes, for more discussion see [5]. Then from eq. (21) we find

$$A_{\pm\mp} \propto \Delta^2 \frac{\alpha_S(Q_V^2)}{(Q^2+m_V^2)^2} \left[\frac{6G(x,Q_V^2)}{Q^2+m_V^2} \frac{\langle k_z^2 \rangle}{4m_V^2} + \frac{G(x,\mu_G^2)}{\mu_G^2} \right], \tag{30}$$

where the leading $\log\overline{Q}^2$ amplitude is of higher twist. Such a mismatch of the twist and leading $\log\overline{Q}^2$ regime in diffractive DIS is dejà vu: the leading twist σ_T is soft-gluon dominated whereas the full fledged $\log\overline{Q}^2$ is at work for the higher twist σ_L [10, 14]. What is new in (30) is that the both regimes mix in one and the same helicity amplitude. The soft-gluon exchange dominance of the leading twist double-helicity flip was noticed recently by Ivanov and Kirschner [15].

Finally, we emphasize that non-vanishing single- and double-helicity flip amplitudes (27) and (29) do not require the applicability of pQCD and can best be searched for in real photo- or electroproduction at small $Q^2 \lesssim m_V^2$.

When this manuscript was under preparation, we learned of the related work by Ivanov and Kirschner (IK) [15]. IK considered only light quarkonia and put $m_q = 0$, our results are applicable from light to heavy quarkonia. While we agree with IK on the twist of helicity amplitudes, he have differences in the form if $I(\gamma^* \to V)$. For instance, the $\gamma^* \leftrightarrow V$ symmetry is not manifest in the IK formulas. The apparent source of differences is the unwarranted omission by IK of terms $\propto \Delta$ in the polarization vectors of vector mesons.

To summarize, we presented the perturbative QCD derivation of helicity amplitudes for diffractive electroproduction of vector mesons. Compared to IK [15] our derivation holds for light to heavy vector mesons and our formulas are applicable also to production of radially excited states. We determined the twist of the s-channel helicity non-conserving amplitudes. With the exception of double-flip, all helicity amplitudes are proportional to the gluon structure function of the proton at a similar pQCD hardness scale (24). Our principal conclusions on substantial s-channel helicity non-conserving effects hold beyond the perturbative QCD and are applicable also to real photoproduction.

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