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Diffractive S and D–wave vector mesons in deep inelastic scattering

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Abstract

We derive helicity amplitudes for diffractive leptoproduction of the S and D wave states of vector mesons. We predict a dramatically different spin dependence for production of the S and D wave vector mesons. We find very small $R = \sigma_L/\sigma_T$ and abnormally large higher twist effects in production of longitudinally polarized D-wave vector mesons.

Diffractive vector meson production $\gamma^* + p \to V + p'$, 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, 2, 3, 4, 5], for the recent review see [6]). (For the kinematics see fig. 1, $Q^2 = -q^2$ and $W^2 = (p+q)^2$ are standard DIS variables). The ground state vector mesons, $V = \rho^0, \omega^0, \varphi^0, J/\Psi, \Upsilon$ are usually supposed to be the S-wave spin-triplet $q\bar{q}$ states. However, all the previous theoretical calculations used the $V\bar{q}q$ vertex $\phi_V V_\mu \bar{q}\Gamma_\mu q$ with the simplest choice $\Gamma_\mu = \gamma_\mu$, which corresponds to a certain mixture of the S- and D-wave states, and any discussion of the impact of the D-wave admixture in the literature is missing (here V_μ is the vector meson polarization vector and ϕ_V is the vertex function related to the vector meson LCWF as specified below.).

We report here a derivation of helicity amplitudes for diffractive production of pure S and D-wave $q\bar{q}$ systems for small to moderate momentum transfer $\mathbf{\Delta}$ within the diffraction cone. Understanding production of D-wave states is a topical issue for several reasons. First, the D-wave admixture may affect predictions for the ratio $R = \sigma_L/\sigma_T$, in which there is a persistent departures of theory from the experiment. To this end we recall that the nonperturbative long-range pion-exchange between light quarks and antiquarks [7] is a natural source of S-D

mixing in the ground state ρ^0 and ω^0 mesons. Second, different spin properties of S- and D-wave production may facilitate as yet unresolved D-wave vs. $2S$-wave assignment of the $\rho'(1480)$ and $\rho'(1700)$ and of the $\omega'(1420)$ and $\omega'(1600)$ mesons.

In our analysis we rely heavily upon the derivation [8] of amplitudes of the s-channel helicity conserving (SCHC) and non-conserving (SCHNC) transitions, albeit in slightly different notations. We predict a dramatically different spin dependence for production of the S and D wave states, especially the Q^2 dependence of $R = \sigma_L/\sigma_T$ which derives from the anomalously large higher twist effects in the SCHC amplitude for production of longitudinally polarized vector mesons. Our technique can be readily generalized to higher excited states, 3^- etc, leptoproduction of which is interesting for the fact that they cannot be formed in e^+e^- annihilation.

A typical leading $\log\frac{1}{x}$ ($\mathrm{LL}\frac{1}{x}$) pQCD diagram for vector meson production is shown in Fig. 1. We use the standard Sudakov expansion of all the momenta in the two lightcone vectors

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

such that $q'^2 = p'^2 = 0$ and $s = 2p' \cdot q'$, and the two-dimensional transverse component: $k = zq' + yp' + k_\perp, \quad \kappa = \alpha q' + \beta p' + \kappa_\perp, \quad \Delta = \gamma p' + \delta q' + \Delta_\perp$ (with the exception of $\mathbf{r}$ which is a 3-dimensional vector, see below, hereafter $\mathbf{k}, \boldsymbol{\Delta}, ..$ always stand for 2-dimensional $k_\perp, \Delta_\perp$ etc.). The diffractive helicity amplitudes take the form

$$A^{S,D}_{\lambda_V\lambda_\gamma}(x, Q^2, \boldsymbol{\Delta}) = is\frac{C_F N_c c_V \sqrt{4\pi\alpha_{em}}}{2\pi^2}\int_0^1 \frac{dz}{z(1-z)}\int d^2\mathbf{k}\psi_{S,D}(z,\mathbf{k})$$
$$\int\frac{d^2\kappa}{\kappa^4}\alpha_S(\max\left\{\kappa^2, \mathbf{k}^2 + \overline{Q}^2\right\})I^{S,D}_{\lambda_V\lambda_\gamma}(\gamma^* \to V)\left(1 + i\frac{\pi}{2}\frac{\partial}{\partial\log x}\right)\mathcal{F}(x,\kappa,\boldsymbol{\Delta}), \qquad (1)$$

where $\lambda_V, \lambda_\gamma$ stand for helicities, m is the quark mass, $C_F = \frac{N_c^2-1}{2N_c}$ is the Casimir operator, $N_c = 3$ is the number of colors, $c_V = \frac{1}{\sqrt{2}}, \frac{1}{3\sqrt{2}}, \frac{1}{3}, \frac{2}{3}$ for the $\rho^0, \omega^0, \phi^0, J/\Psi$ mesons, α_{em} is the fine structure constant, α_S is the strong coupling and $\overline{Q}^2 = m^2 + z(1-z)Q^2$ is the relevant hard scale. To the $\mathrm{LL}\frac{1}{x}$ the lower blob is related to the unintegrated gluon density matrix $\mathcal{F}(x,\kappa,\boldsymbol{\Delta})$ [5, 10, 11]. For small $\boldsymbol{\Delta}$ within the diffraction cone

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

where $\partial G/\partial\log\kappa^2$ is the conventional unintegrated gluon structure function and, modulo to a slow Regge growth, the diffraction cone $B_{3\mathbf{I\!P}} \sim 6$ GeV^{-2} [5].

In the light–cone formalism [9], one first computes the production of an on-mass shell $q\bar{q}$ pair of invariant mass M and total momentum q_M. This amplitude is projected onto the state $(q\bar{q})_J$ of total angular momentum $J = 1$ using the running longitudinal and the usual transverse polarization vectors

$$V_L = \frac{1}{M}\left(q' + \frac{\boldsymbol{\Delta}^2 - M^2}{s}p' + \Delta_\perp\right), \qquad V_T = V_\perp + \frac{2(\mathbf{V}_\perp\cdot\boldsymbol{\Delta})}{s}(p' - q'), \qquad (3)$$

such that $(V_T V_L) = (V_T q_M) = (V_L q_M) = 0$. Then the resulting upper blob $I(\gamma^* \to V)$ is contracted with the radial LCWF of the $q\bar{q}$ Fock state of the vector meson,

$$\psi_{S,D}(z,\mathbf{k}) = \psi_{S,D}(\mathbf{r}^2) = \frac{\phi_{S,D}(\mathbf{r}^2)}{M^2 - m_V^2}\,. \qquad (4)$$

Here $r = \frac{1}{2}(k_2 - k_1)$, which in the rest frame is the relative 3-momentum in the $q\bar{q}$ pair, $r = (0, \mathbf{r}) = (0, \mathbf{k}, k_z)$, $r^2 = -\mathbf{r}^2$, and

$$M^2 = 4(m^2 + \mathbf{r}^2) = \frac{m^2 + \mathbf{k}^2}{z(1-z)}\,.$$

To conform to this procedure, all the occurrences of the vector meson mass m_V in $I_{\lambda_V \lambda_\gamma}$ of ref. [8] must be replaced by M.

A useful normalization of the radial LCWF's $\psi_{S,D}(\mathbf{r}^2)$ is provided by the $V \to e^+e^-$ decay constant, $\langle 0 | J^{em}_{mu} | V \rangle = f c_V \sqrt{4\pi\alpha_{em}} V_\mu$:

$$f_S = \frac{N_c}{(2\pi)^3} \int d^3\mathbf{r} \frac{8}{3}(M+m)\psi_S(\mathbf{r}^2)\,, \quad f_D = \frac{N_c}{(2\pi)^3} \int d^3\mathbf{r} \frac{32}{3} \frac{\mathbf{r}^4}{M+2m} \psi_D(\mathbf{r}^2)\,. \qquad (5)$$

The nice observation is that we need not go again through all the calculations of helicity amplitudes. Indeed, the spinor vertices $\Gamma^{S,D}_\mu$ for the pure S and D wave states can be readily obtained from the simplest $\Gamma_\mu = \gamma_\mu$ used in [8]. Following [9], it can be easily shown that

$$\Gamma^S_\mu = \gamma_\mu - \frac{2r_\mu}{M+2m} = \mathcal{S}_{\mu\nu}\gamma_\nu; \quad \mathcal{S}_{\mu\nu} = g_{\mu\nu} - \frac{2r_\mu r_\nu}{m(M+2m)}. \qquad (6)$$

Here we made use of $r^\mu \gamma_\mu = m$ and $(q_M \cdot r) = 0$. Once the S-wave is constructed, the spinor structure for a D-wave state can be readily obtained by contracting the S-wave vertex with $3r_\mu r_\nu + g_{\mu\nu}\mathbf{r}^2$ with the result

$$\Gamma^D_\mu = \mathbf{r}^2\gamma_\mu + (M+m)r_\mu = \mathcal{D}_{\mu\nu}\gamma_\nu; \quad \mathcal{D}_{\mu\nu} = \mathbf{r}^2 g_{\mu\nu} + \frac{M+m}{m} r_\mu r_\nu. \qquad (7)$$

Consequently, the answers for either S or D-wave production amplitudes can be immediately obtained from the expressions given in [8] by substitutions $V^*_\mu \to V^*_\nu \mathcal{S}_{\nu\mu}$, $V^*_\mu \to V^*_\nu \mathcal{D}_{\nu\mu}$ for S and D-wave states respectively.

In terms of diffractive amplitudes $\mathbf{\Phi}_1$ and Φ_2 defined in [8], we find for S-wave vector mesons (here T stands for the transverse polarization)

$$I^S_{0L} = -4QMz^2(1-z)^2\Phi_2 \left[1 + \frac{(1-2z)^2 m}{2z(1-z)(M+2m)}\right]\,,$$

$$I^S_{TT} = \Big\{(\mathbf{V}^*\mathbf{e})[m^2\Phi_2 + (\mathbf{k}\mathbf{\Phi}_1)] + (1-2z)^2(\mathbf{V}^*\mathbf{k})(\mathbf{e}\mathbf{\Phi}_1)\frac{M}{M+2m}$$
$$-(\mathbf{e}\mathbf{k})(\mathbf{V}^*\mathbf{\Phi}_1) + \frac{2m}{M+2m}(\mathbf{V}^*\mathbf{k})(\mathbf{e}\mathbf{k})\Phi_2\Big\}\,,$$

$$I^S_{0T} = -2z(1-z)(2z-1)M(\mathbf{e}\mathbf{\Phi}_1)\left[1 + \frac{(1-2z)^2 m}{2z(1-z)(M+2m)}\right] + \frac{Mm}{M+2m}(2z-1)(\mathbf{e}\mathbf{k})\Phi_2\,,$$

$$I^S_{TL} = 2Qz(1-z)(2z-1)(\mathbf{V}^*\mathbf{k})\Phi_2 \frac{M}{M+2m}\,. \qquad (8)$$

Because the difference between Γ^S_μ and γ_μ is a relativistic correction, the results for the S-wave vector mesons differ from those found in [8] only by a small relativistic corrections $\propto \mathbf{r}^2/M^2$. The exceptional case is suppression of I^S_{TL} by factor $M/(2m+M) \sim 0.5$.

We skip the twist expansion for S-wave amplitides, which can easily be done following [8], and proceed to the much more interesting case of D-wave mesons, for which

$$I^D_{0L} = -QMz(1-z)\left(\mathbf{k}^2 - \frac{4m}{M}k_z^2\right)\Phi_2\,,$$
$$I^D_{TT} = \Big\{(\mathbf{V}^*\mathbf{e})\mathbf{r}^2[m^2\Phi_2 + (\mathbf{k}\mathbf{\Phi}_1)] + (1-2z)^2(\mathbf{r}^2+m^2+Mm)(\mathbf{V}^*\mathbf{k})(\mathbf{e}\mathbf{\Phi}_1)$$
$$-\mathbf{r}^2(\mathbf{e}\mathbf{k})(\mathbf{V}^*\mathbf{\Phi}_1) - m(M+m)(\mathbf{V}^*\mathbf{k})(\mathbf{e}\mathbf{k})\Phi_2\Big\}\,,$$
$$I^D_{0T} = -\frac{2z-1}{2}M\left\{(\mathbf{e}\mathbf{\Phi}_1)(\mathbf{k}^2 - \frac{4m}{M}k_z^2) + m(M+m)(\mathbf{e}\mathbf{k})\Phi_2\right\}\,,$$
$$I^D_{TL} = 2Qz(1-z)(2z-1)(\mathbf{V}^*\mathbf{k})(\mathbf{r}^2+m^2+Mm)\Phi_2\,, \qquad (9)$$

The novel features of these amplitudes are best seen in the twist expansion in inverse powers of the hard scale $\overline{Q}^2$. As it was noted in [8], 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 $\mathbf{k}^2 \sim R_V^{-2}, \mathbf{\Delta}^2 \ll \kappa^2 \ll \overline{Q}^2$. The closer inspection of our $I^D_{\lambda_V\lambda_\gamma}$ shows that the seemingly leading interference with the dominant S-wave component in the photon always appears in the quadrupole combination $2k_z^2 - \mathbf{k}^2$. Since the integration over quark loop can be cast in form $d^3\mathbf{r}$, such quadrupole combinations vanish after angular integration. As a result, the abnormally large higher twist contributions $\propto M^2/(M^2+Q^2$ with large numerical factors come into play and significantly modify the Q^2 dependence of amplitudes for production of longitudinally polarized vector mesons:

$$I^D_{0L} = -\frac{Q}{M}\cdot\frac{32\mathbf{r}^4}{15(M^2+Q^2)^2}\cdot\left(1 - 8\frac{M^2}{M^2+Q^2}\right)\kappa^2\,, \qquad (10)$$
$$I^D_{\pm\pm} = (\mathbf{V}^*\mathbf{e})\cdot\frac{32\mathbf{r}^4}{15(M^2+Q^2)^2}\cdot\left(15 + 4\frac{M^2}{M^2+Q^2}\right)\kappa^2\,, \qquad (11)$$
$$I^D_{\pm L} = -\frac{32\mathbf{r}^4}{15(M^2+Q^2)^2}\cdot\frac{24Q(\mathbf{V}^*\mathbf{\Delta})}{M^2+Q^2}\kappa^2\,, \qquad (12)$$
$$I^D_{L\pm} = \frac{32\mathbf{r}^4}{15(M^2+Q^2)^2}\cdot\frac{8(\mathbf{e}\mathbf{\Delta})}{M}\left(1 + 3\frac{M^2}{M^2+Q^2}\right)\kappa^2\,, \qquad (13)$$
$$I^D_{\pm\mp} = (\mathbf{V}^*\mathbf{\Delta})(\mathbf{e}\mathbf{\Delta})\cdot\frac{32\mathbf{r}^4}{15(M^2+Q^2)^2}\cdot\left(1 - \frac{96}{7}\frac{\kappa^2\mathbf{r}^2}{M^2(M^2+Q^2)}\right)\,. \qquad (14)$$

In a close similarity to the S-wave case [8], the leading twist double-helicity flip amplitude is dominated by soft gluon exchange, the LL$\overline{Q}^2$ component is of higher twist.

In order to emphasize striking difference between the D-wave and S-wave state amplitudes, we focus on nonrelativistic heavy quarkonia, where $M^2 \approx m_V^2$, although all the qualitative results hold for light vector mesons too. For the illustration purposes, we evaluated the ratios of helicity amplitudes, $\rho_{D/S} = f_S A^D/f_D A^S$, for the the harmonic oscillator wave functions:

$$\rho_{0L}(D/S) = \frac{1}{5}\left(1 - 8\frac{m_V^2}{Q^2+m_V^2}\right)\,,$$
$$\rho_{\pm\pm}(D/S) = 3\left(1 + \frac{4}{15}\frac{m_V^2}{Q^2+m_V^2}\right)\,,$$

$$\rho_{0\pm}(D/S) = -\frac{1}{5}(m_V a_S)^2\left(1 + 3\frac{m_V^2}{Q^2+m_V^2}\right)\,,$$
$$\rho_{\pm L}(D/S) = \frac{3}{40}(m_V a_S)^4\,. \qquad (15)$$

First, A_{0L} changes the sign at $Q^2 \sim 7m_V^2$. The ratio $R^D = \sigma_L/\sigma_T$ has thus a non-monotonous Q^2 behavior and $R^D \ll R^S$. Furthermore, $R^D \lesssim 1$ in a broad range of $Q^2 \lesssim 225m_V^2$. Whereas in heavy quarkonia the S–D mixing is arguably weak [12], in light ρ^0, ω^0 even a relatively weak S-D mixing could have a substantial impact on R. Second, all the D-wave amplitides, SCHC and SCHNC alike, with exception of the higher twist component of double-helicity flip, are proportional to $\mathbf{r}^4$ and, in view of eq. (5), to the decay constant f_D. In contrast to that, in the S-wave case the spin-flip amplitudes for heavy quarkonia are suppressed by nonrelativistic Fermi motion [8]. The relevant suppression parameter is $\sim 1/(a_S m_V)^2$, where a_S is the radius of the $1S$ state. For this reason, for D-wave states the SCHNC effects are much stronger. For instance, for the charmonium $(m_V a_S)^2 \approx 27$, see [12].

To summarize, we found dramatically different spin properties of diffractive leptoproduction of the S and D wave states of vector mesons. We predict very small $R^D = \sigma_L/\sigma_T$ and very strong breaking of s-channel helicity conservation in production of D-wave states. Higher twist effects in production of longitudinally polarized D-wave vector mesons are found to be abnormally large. Consequently, the distinct spin properties of D-wave vector mesons in diffractive DIS offer an interesting way to discern S and D-wave meson states, which are indistinguishable at e^+e^- colliders.

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Figure caption:

Fig.1: One of the four Feynman diagrams for the vector meson production $\gamma^* p \to V p'$ via QCD two-gluon pomeron exchange.