Remarks on exclusive electroproduction of transversely polarized vector mesons

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We discuss the electroproduction of light vector mesons from transversely polarized photons. Here QCD factorization cannot be applied as shown explicitly in a leading-order calculation. It is emphasized that present infrared singular contributions cannot be regularized through phenomenological meson distribution amplitudes with suppressed end-point configurations. We point out that infrared divergences arise also from integrals over skewed nucleon parton distributions. In a phenomenological analysis of transverse vector meson production model-dependent regularizations have to be applied. If this procedure preserves the structure suggested by a leading-order calculation of Feynman diagrams, one obtains contributions from nucleon parton distributions and their derivatives. In particular, polarized gluons enter only through their derivative.

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Exclusive electroproduction of mesons from nucleons has become a topic of broad interest. Experimental and theoretical advances have supported this development. At high energies a large amount of data has become available from experiments at CERN [New Muon Collaboration (NMC)] and DESY (HERA, HERMES) (for references see, e.g., [1,2]). Further measurements were carried out at DESY, and discussed at CERN (COMPASS) [3,4] and TJNAF [5]. From the theoretical side a factorization theorem proven in Ref. [6] set the basis for many investigations. It states that the underlying photon-parton subprocesses are dominated for longitudinally polarized photons and large photon virtualities, $Q^2 \gg \Lambda^2_{\rm QCD}$, by short distances and, hence, can be calculated perturbatively.

As already emphasized in [6], the interaction of transversely polarized photons cannot be treated in a framework based on QCD factorization. This is due to infrared sensitive contributions which, e.g., result from large size quark-antiquark configurations in the produced meson. A straightforward QCD analysis of transverse vector meson production, as done for longitudinal ones (see Refs. [7–13]), is therefore not possible. On the other hand, expected new data will provide more detailed information on these processes and help to investigate strong interaction dynamics at large distances.

Descriptions of exclusive meson production from transversely polarized photons rely on model assumptions which are needed to regularize the present, at least logarithmic, infrared singularities. The exploratory analysis presented in this note is based on a naive application of Wilson's operator product expansion. In this framework we find constraints on the vector meson production amplitude imposed by Lorentz covariance. Different vector meson distribution amplitudes, which enter the production amplitude, are related by Wandura-Wilczek-type relations [14,15]. These have significant implications for the nature of the present infrared singularities.

In particular we observe that infrared divergences arise from integrals over meson distribution amplitudes *and* from integrals over skewed nucleon parton distributions. Furthermore, our calculation yields an analytic structure of the production amplitude for transverse vector mesons which is richer than for longitudinal ones. In particular, we obtain contributions which involve both parton distributions and their derivatives.

In the following we restrict ourselves to the helicity conserving, hard exclusive production of ρ mesons from transversely polarized photons. An important ingredient of the corresponding amplitude are meson distribution amplitudes. Following Refs. [14,15] we parametrize the meson-to-vacuum matrix elements of vector and axial vector[1] currents as

$$\begin{aligned}\langle 0|\overline{q}(0)[0;x]\gamma_\mu q(x)|\rho(p,\lambda)\rangle_{x^2=0}\\ = p_\mu \frac{\mathcal{E}^{(\lambda)}\cdot x}{p\cdot x} f_\rho m_\rho \int_0^1 d\tau\, e^{-i\tau p\cdot x}\phi_{||}(\tau)\\ + \left(\mathcal{E}^{(\lambda)}_\mu - p_\mu \frac{\mathcal{E}^{(\lambda)}\cdot x}{p\cdot x}\right) f_\rho m_\rho \int_0^1 d\tau\, e^{-i\tau p\cdot x} g^{v}(\tau),\end{aligned} \quad (1)$$

and

$$\begin{aligned}\langle 0|\overline{q}(0)[0;x]\gamma_\mu\gamma_5 q(x)|\rho(p,\lambda)\rangle_{x^2=0}\\ = \frac{1}{4}\varepsilon_{\mu\nu\rho\sigma}\mathcal{E}^{(\lambda)\nu}p^\rho x^\sigma f_\rho m_\rho \int_0^1 d\tau\, e^{-i\tau p\cdot x} g^{a}(\tau).\end{aligned} \quad (2)$$

Here p^μ is the four momentum of the ρ meson with invariant mass m_ρ and decay constant f_ρ. The vector meson polarization is specified by λ which corresponds to the polarization vector $\mathcal{E}^{(\lambda)}$. The light-cone matrix elements, as well as the distribution amplitudes in Eqs. (1), (2), are defined at a cer-

[1] We use the convention of [16] for γ_5 and the epsilon tensor.

tain renormalization scale μ which we suppress if convenient. Gauge invariance is guaranteed by the path-ordered exponential

$$[0;x]=\mathcal{P}\exp\left[igx_\mu\int_0^1 A^\mu(x\eta)d\eta\right],$$

which reduces to 1 in axial gauge $n\cdot A=0$ (g stands for the strong-coupling constant and A^μ denotes the gluon field). In light-cone gauge the twist-2 distribution amplitude $\phi_{||}(\tau)$ can be related to the wave function of the minimal quark-antiquark Fock state of zero orbital angular momentum in a longitudinally polarized meson [17–20]. Note that the twist-2 distribution amplitude for a transversally polarized ρ, related to the zero angular-momentum component of its quark-antiquark wave function, is determined by the matrix element of a chiral-odd tensor quark operator and does not contribute to the process considered here.

In the following we use the antisymmetric distribution $\Phi_{||}(\tau)$ given by [14]

$$\Phi_{||}(\tau)=\frac{1}{2}\left[\bar{\tau}\int_0^\tau du\,\frac{\phi_{||}(u)}{\bar{u}}-\tau\int_\tau^1 du\,\frac{\phi_{||}(u)}{u}\right], \tag{3}$$

with $\bar{\tau}=1-\tau$. As explained in [14,15] twist-2 and twist-3 string operators contribute to the distributions $g^v(\tau)$ and $g^a(\tau)$. Lorentz covariance leads to the presence of twist-2 contributions given by Wandzura-Wilczek-type relations:

$$g^v(\tau)=\frac{1}{2}\left[\int_0^\tau du\,\frac{\phi_{||}(u)}{\bar{u}}+\int_\tau^1 du\,\frac{\phi_{||}(u)}{u}\right], \tag{4}$$

$$g^a(\tau)=2\left[\bar{\tau}\int_0^\tau du\,\frac{\phi_{||}(u)}{\bar{u}}+\tau\int_\tau^1 du\,\frac{\phi_{||}(u)}{u}\right]. \tag{5}$$

Both, $g^v(\tau)$ and $g^a(\tau)$ are symmetric functions of τ. Twist-3 contributions are related to matrix elements of three-particle quark-gluon-quark operators [15] and will not be considered here. The QCD evolution of three-particle distribution amplitudes is different from that of two-particle ones. Therefore, their contribution to the end-point behavior of $g^v(\tau)$ and $g^a(\tau)$ can be discussed separately since a cancellation of infrared divergences can occur only accidentally, i.e., only for one particular renormalization scale.

In general, one can assume $\sigma_\phi=\int_0^1 du\,\phi_{||}(u)/u$ is different from zero. This is quite natural since $\sigma_\phi=0$ can be fulfilled, if at all, only at one particular scale μ due to the scale dependence of $\phi_{||}(u;\mu)$ [14]. As a consequence one finds from Eqs. (4) and (5) in the limit $\tau\to 0$: $g^v(\tau)\sim\sigma_\phi$, $g^a(\tau)\sim\sigma_\phi\tau$, and $\Phi_{||}(\tau)\sim\sigma_\phi\tau$.

The amplitude $\mathcal{M}^{\gamma_\perp^*\to\rho_\perp}$ for ρ meson production from transversally polarized virtual photons can be split into parts

$$\mathcal{M}^{\gamma_\perp^*\to\rho_\perp}=\mathcal{M}_G^{\gamma_\perp^*\to\rho_\perp}+\mathcal{M}_q^{\gamma_\perp^*\to\rho_\perp}+\mathcal{M}_{\Delta G}^{\gamma_\perp^*\to\rho_\perp}+\mathcal{M}_{\Delta q}^{\gamma_\perp^*\to\rho_\perp}, \tag{6}$$

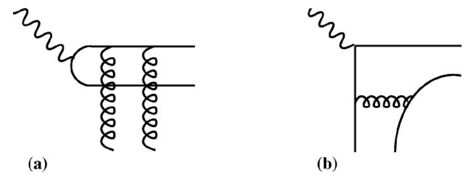

FIG. 1. Typical contributions to hard exclusive meson production.

involving unpolarized and polarized skewed quark and gluon distribution functions, respectively. A calculation of leading-order Feynman diagrams can be done along the lines of Ref. [11].

For illustration we show intermediate steps for the gluon contribution. In second-order perturbation theory the S matrix reads

$$S=i\frac{g^2}{2}\int d^4x\,d^4y\,d^4z\,e^{-iq\cdot x}\times\langle\rho(q',\lambda)N(P',S')|\mathcal{M}(x,y,z)|N(P,S)\rangle, \tag{7}$$

with the time-ordered product of quark and gluon fields

$$\mathcal{M}(x,y,z)=T[\bar{\psi}^a(z)\gamma^\rho\psi^b(z)\bar{\psi}^c(x)\not{\varepsilon}_T\psi^c(x)\bar{\psi}^d(y)\times\gamma^\sigma\psi^e(y)A_\rho^A(z)A_\sigma^B(y)]t_{ab}^A t_{de}^B. \tag{8}$$

The incident photon carries a four-momentum q and is transversely polarized as specified by the polarization vector ε_T. The generators of color SU(3) are denoted by t^A and t^B. The momenta and spins of the initial and scattered nucleon are P,P' and S, S', respectively; the produced meson carries momentum q' and polarization λ.

As a next step we perform the usual operator product expansion of $\mathcal{M}(x,y,z)$. Choosing, for example, the flow of the hard momentum q as in Fig. 1(a) one obtains a term which can be interpreted as photon-meson transition in a background gluon field provided by the nucleon. A different choice of the hard momentum flow results in a photon-meson transition in the background of quark fields as illustrated in Fig. 1(b). In other, not shown, leading-order contributions the photon couples to one of the other possible quark lines in Fig. 1.

For the diagram in Fig. 1(a) we find in leading-order perturbation theory:

$$\mathcal{M}(x,y,z)\ni\bar{\psi}^a(z)\gamma^\rho S_q^{(bd)}(z-y)\gamma^\sigma S_q^{(ec)}(y-x)\times\not{\varepsilon}_T\bar{\psi}^c(x)A_\rho^A(z)A_\sigma^B(y)t_{ab}^A t_{de}^B, \tag{9}$$

where S_q stands for the perturbative quark propagator. As a next step we perform a Fierz transformation in color and Dirac space, and project onto color singlet pieces. We then obtain terms with bilinear quark operators which transform as vector and axial vector, respectively. For simplicity we restrict ourselves to the axial vector piece. It reads

$$\mathcal{M}(x,y,z) \ni \frac{\delta^{AB}}{8N_c}\overline{\psi}(z)\gamma_\mu\gamma_5\psi(x)\mathrm{Tr}[\gamma^\rho S_q(z-y)\gamma^\sigma \times S_q(y-x)\not{\varepsilon}_T\gamma^5\gamma^\mu]A^A_\rho(z)A^B_\sigma(y). \tag{10}$$

After taking the matrix element (7), the meson distribution amplitudes from Eq. (2) appear. The skewed unpolarized and polarized gluon distributions of the nucleon, $G(u,\xi)$ and $\Delta G(u,\xi)$, enter through the nucleon matrix element of the gluon operators in Eq. (10). In the light-cone gauge they are defined as [9,21]

$$\begin{aligned}\langle N(P',S')|A^A_\rho(0)A^A_\sigma(y)|N(P,S)\rangle_{y^2=0} &= \frac{1}{4}(-g^T_{\rho\sigma})\frac{\overline{N}(P',S')\not{n}N(P,S)}{\overline{P}\cdot n}e^{-r/2\cdot y}\int_{-1}^1 du\frac{G(u,\xi)}{(u-\xi+i\epsilon)(u+\xi-i\epsilon)}e^{-iu\overline{P}\cdot y}\\ &+\frac{1}{4}i\varepsilon_{\rho\sigma\alpha\beta}\frac{n^\alpha n^{*\beta}}{n\cdot n^*}\frac{\overline{N}(P',S')\gamma_5\not{n}N(P,S)}{\overline{P}\cdot n}e^{-r/2\cdot y}\int_{-1}^1 du\frac{\Delta G(u,\xi)}{(u-\xi+i\epsilon)(u+\xi-i\epsilon)}e^{-iu\overline{P}\cdot y}.\end{aligned} \tag{11}$$

$N(P,S)$ and $\overline{N}(P',S')$ are the Dirac spinors of the initial and scattered nucleon, respectively. The average nucleon momentum is $\overline{P}=(P+P')/2$, and the momentum transfer is $r=P-P'$. Furthermore, we have introduced the momentum $\overline{q}=(q+q')/2$ with $\overline{Q}^2=-\overline{q}^2$, $\overline{\omega}=2\overline{q}\cdot\overline{P}/(-\overline{q}^2)$ and $\xi=1/\overline{\omega}$. Finally, n and n^* are lightlike vectors with $n\cdot a=a^+=a^0+a^3$ and $n^*\cdot a=a^-=a^0-a^3$ for any vector a. We then define $-g^T_{\rho\sigma}=(n^*_\rho n_\sigma+n_\rho n^*_\sigma)/n\cdot n^*-g_{\rho\sigma}$. In the forward limit, $\xi\to 0$, the skewed gluon distribution functions reduce to the ordinary unpolarized and polarized gluon distributions of the nucleon:

$$\begin{aligned}\lim_{\xi\to 0}G(u,\xi)&=ug(u),\\ \lim_{\xi\to 0}\Delta G(u,\xi)&=u\Delta g(u).\end{aligned} \tag{12}$$

For simplicity we have omitted so-called K terms, which arise in the matrix element (11) [9,21]. Their contribution to the production amplitudes (15), (16) can be obtained simply by replacing the skewed gluon distributions by K distributions (including corresponding Dirac prefactors) [11].

Performing the integrations in Eq. (7) one obtains from Eq. (10):

$$\begin{aligned}\mathcal{M}^{\gamma^*_\perp\to\rho_\perp}_{\Delta G}+\mathcal{M}^{\gamma^*_\perp\to\rho_\perp}_{G} &\ni \frac{g^2}{256N_c}\frac{f_\rho m_\rho}{\overline{P}\cdot n}\epsilon_{\mu\nu\delta\gamma}\mathcal{E}^{*\nu}_T q'^\delta\int_0^1 d\tau\, g^a(\tau)\int_{-1}^1 du\int d^4k\\ &\times\frac{\mathrm{Tr}\{\gamma^\rho[\not{k}+\overline{\not{P}}(\xi+u)]\gamma^\sigma\not{k}\not{\varepsilon}_T\gamma_5\gamma^\mu\}}{\{[k+\overline{P}(\xi+u)]^2+i\epsilon\}(k^2+i\epsilon)(u-\xi+i\epsilon)(u+\xi-i\epsilon)}\frac{\partial}{\partial k_\gamma}\delta^4(\tau q'-2\xi\overline{P}-k)\\ &\times\overline{N}(P',S')\left\{g^T_{\rho\sigma}\not{n}G(u,\xi)+i\epsilon_{\rho\sigma\alpha\beta}\frac{n^\alpha n^{*\beta}}{n\cdot n^*}\gamma_5\not{n}\Delta G(u,\xi)\right\}N(P,S).\end{aligned} \tag{13}$$

Integrating by parts gives

$$\begin{aligned}\mathcal{M}^{\gamma^*_\perp\to\rho_\perp}_{\Delta G}+\mathcal{M}^{\gamma^*_\perp\to\rho_\perp}_{G} &\ni -i\frac{g^2}{128N_c}\frac{f_\rho m_\rho}{\overline{Q}^2}\int_0^1\frac{d\tau}{\tau^2}g^a(\tau)\left\{\varepsilon_T\cdot\mathcal{E}^*_T\frac{\overline{N}(P',S')\not{n}N(P,S)}{\overline{P}\cdot n}\int_{-1}^1 du\frac{G(u,\xi)}{\xi-u-i\varepsilon}\left[\frac{1}{\xi-u-i\varepsilon}+\frac{1}{\xi+u-i\varepsilon}\right]\right.\\ &\left.-i\epsilon_{\mu\nu\alpha\beta}\frac{n^\alpha n^{*\beta}}{n\cdot n^*}\varepsilon^\mu_T\mathcal{E}^{*\nu}_T\frac{\overline{N}(P',S')\gamma_5\not{n}N(P,S)}{\overline{P}\cdot n}\int_{-1}^1 du\frac{\Delta G(u,\xi)}{\xi-u-i\varepsilon}\left[\frac{u\overline{\omega}}{\xi-u-i\varepsilon}+\frac{u\overline{\omega}}{\xi+u-i\varepsilon}\right]\right\}.\end{aligned} \tag{14}$$

The remaining leading-order contributions to the production amplitude $\mathcal{M}^{\gamma^*_\perp\to\rho_\perp}$ can be calculated in a similar way. In total one obtains for the amplitude which involves the unpolarized gluon distribution:

$$\mathcal{M}_G^{\gamma_\perp^*\to\rho_\perp}=i\frac{g^2}{32N_C}\frac{f_\rho m_\rho}{\bar{Q}^2}\frac{\bar{N}(P',S')\not{n}N(P,S)}{\bar{P}\cdot n}\varepsilon_T\cdot\mathcal{E}_T^*\Bigg\{\mathcal{I}_1\int_{-1}^1 du\, G(u,\xi)\left[\frac{\bar{\omega}}{\xi-u-i\epsilon}+\frac{\bar{\omega}}{\xi+u-i\epsilon}\right]$$
$$+\mathcal{I}_2\int_{-1}^1 du\, G(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)^2}+\frac{1}{(\xi+u-i\epsilon)^2}\right]\Bigg\}, \tag{15}$$

while the contribution from polarized gluons reads

$$\mathcal{M}_{\Delta G}^{\gamma_\perp^*\to\rho_\perp}=-\frac{g^2}{32N_C}\frac{f_\rho m_\rho}{\bar{Q}^2}\frac{\bar{N}(P',S')\gamma_5\not{n}N(P,S)}{\bar{P}\cdot n}\times\varepsilon_{\mu\nu\alpha\beta}\frac{n^\alpha n^{*\beta}}{n\cdot n^*}\varepsilon_T^\mu\mathcal{E}^{\nu*}\mathcal{I}_2\int_{-1}^1 du\,\Delta G(u,\xi)\times\left[\frac{1}{(\xi-u-i\epsilon)^2}-\frac{1}{(\xi+u-i\epsilon)^2}\right]. \tag{16}$$

The integrals $\mathcal{I}_1$ and $\mathcal{I}_2$ contain the dependence of the production amplitudes on the meson distributions g^a, g^v, and $\Phi_{||}$:

$$\mathcal{I}_1=\int_0^1\frac{d\tau}{\tau}\left(4g^v(\tau)-2\frac{\Phi_{||}(\tau)}{\tau}+\frac{g^a(\tau)}{2\tau\bar{\tau}}\right),$$
$$\mathcal{I}_2=\int_0^1\frac{d\tau}{\tau}\left(2g^v(\tau)+\frac{g^a(\tau)}{2\tau\bar{\tau}}\right). \tag{17}$$

Both, $\mathcal{I}_1$ and $\mathcal{I}_2$ are divergent due to the behavior of the integrands at $\tau\to 0$. These infrared divergences make QCD factorization impossible [6,9]. Currently no QCD framework is available to deal with this problem of infrared sensitive contributions. As a consequence all predictions for ρ production from transversely polarized photons are model dependent. One of the points we want to emphasize in this paper is, that a modification of the end-point behavior of the twist-2 amplitude $\phi_{||}(\tau)$ cannot cure the infrared divergence in Eq. (17). This is a direct consequence of the end-point behavior of $g^v(\tau)$ and $g^a(\tau)$ which is enforced by the Ball-Braun relations (4) and (5).

For the asymptotic distribution amplitude $\phi_{||}(\tau)=6\tau\bar{\tau}$ the integrands in Eq. (17) are proportional to $12/\tau$ and $9/\tau$, respectively. One, therefore, might expect that any phenomenological regularization which is applied to render the integrals finite (see, e.g., [22,23]) leads to a ratio $\mathcal{I}_1/\mathcal{I}_2$ close to one.

Another important point is the dependence of the production amplitudes on the unpolarized and polarized gluon distributions. According to Eq. (15) the unpolarized skewed gluon distribution enters $\mathcal{M}_G^{\gamma_\perp^*\to\rho_\perp}$ through

$$\int_{-1}^1 du\, G(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)}+\frac{1}{(\xi+u-i\epsilon)}\right], \tag{18}$$

and

$$\int_{-1}^1 du\, G(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)^2}+\frac{1}{(\xi+u-i\epsilon)^2}\right]. \tag{19}$$

The first integral is also present in the leading twist production amplitude of ρ mesons via longitudinally polarized photons (see, e.g., [11]). The second contribution, which involves the square of $(\xi\pm u-i\epsilon)$ in the denominator, has not been considered before [22]. We believe that any model of transverse ρ production should include both contributions. In deeply virtual Compton scattering (DVCS) similar contributions have been discussed recently in [24].) Note that the integrals (18) and (19) are well defined only if $G(u,\xi)$ and its first derivative are continuous at $u=\xi$. Although little is known about properties of skewed parton distributions from first principles, at least the asymptotic distribution $G(u,\xi;\mu\to\infty)$ fulfills this requirements [25].

The polarized gluon distribution enters the production amplitude (16) via the integral

$$\int_{-1}^1 du\,\Delta G(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)^2}-\frac{1}{(\xi+u-i\epsilon)^2}\right]. \tag{20}$$

Also this integral exists only if $\Delta G(u,\xi)$ and its derivative are continuous at $u=\xi$ which, again, is suggested by the asymptotic solution of the corresponding QCD evolution equation.

The dependence on the skewed gluon distribution in Eq. (20) is identical to the one found in J/Ψ production [26], but at variance with the model proposed in [23]. Integrating Eq. (20) by parts shows that $\mathcal{M}_{\Delta G}^{\gamma_\perp^*\to\rho_\perp}$ depends on the derivative of $\Delta G(u,\xi)$ rather then on $\Delta G(u,\xi)$ itself. The imaginary part of $\mathcal{M}_{\Delta G}^{\gamma_\perp^*\to\rho_\perp}$ is, for example, proportional to

$$\mathrm{Im}\int_{-1}^1 du\,\Delta G(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)^2}-\frac{1}{(\xi+u-i\epsilon)^2}\right]=-2\pi\frac{\partial}{\partial u}\Delta G(u,\xi)\Big|_{u=\xi}. \tag{21}$$

It can be argued that $\Delta G(u,\xi)$ is proportional to $u\Delta g(u)$ for small ξ and $u\gtrsim\xi$ [11,27]. The magnitude of the imaginary part (21) in the small-ξ region depends then crucially on the small-u behavior of $\Delta g(u)$. Assuming $\Delta g(u)\sim u^{-\lambda}$ for small u leads to an increase of the contribution of polarized gluons to ρ production by a factor $(1-\lambda)/u$ as compared to the model calculation in Ref. [23].

So far we have considered only the gluonic part of the production amplitude (6). The quark amplitudes $\mathcal{M}_q^{\gamma_\perp^* \to \rho_\perp}$ and $\mathcal{M}_{\Delta q}^{\gamma_\perp^* \to \rho_\perp}$ have a similar form, but involve skewed unpolarized and polarized quark distributions $F(u,\xi)$ and $\Delta F(u,\xi)$. Both enter through integrals involving denominators $(\xi \pm u - i\epsilon)$ and their square as in Eqs. (18), (19), (20). However, for the asymptotic flavor singlet quark distribution [25]

$$F(u,\xi;\mu\to\infty) = \frac{15}{2}\,\frac{N_F}{4C_F+N_F}\,\frac{1}{\xi^2}\,\Theta(\xi-|u|)\frac{u}{\xi}\left[1-\left(\frac{u}{\xi}\right)^2\right], \tag{22}$$

one finds an important difference as compared to the gluon case: the real part of the integral

$$\int_{-1}^{1} du\, F(u,\xi)\left[\frac{1}{(\xi-u-i\epsilon)^2}-\frac{1}{(\xi+u-i\epsilon)^2}\right], \tag{23}$$

which is present in $\mathcal{M}_q^{\gamma_\perp^* \to \rho_\perp}$, is infrared singular. $F(u,\xi;\mu)$ approaches its asymptotic form in a continuous way [25], therefore this divergence should be present already at a lower normalization scale μ_0.[2] This observation makes clear that a phenomenological regularization of the integrals which involve meson distribution amplitudes is not sufficient to obtain a finite result for the complete production amplitude.

In summary, we have shown via an explicit calculation of leading-order Feynman diagrams how infrared singular contributions enter in the production of light vector mesons through the interaction of transversely polarized photons. They arise through integrals over involved meson production amplitudes and skewed parton distributions. The former cannot be regularized through phenomenological meson distribution amplitudes with suppressed end-point configurations. This is a direct consequence of Lorentz invariance which provides relations between different vector meson distributions.

In a phenomenological analysis of data on transverse ρ meson production model-dependent regularizations have to be applied. If such a procedure preserves the structure suggested by a leading-order calculation of Feynman diagrams, one obtains amplitudes which involve contributions from parton distribution functions and their derivatives. In particular polarized gluons enter only through their derivative.

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