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Central production of mesons: Exotic states versus Pomeron structure

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Abstract

We demonstrate that the azimuthal dependence of central meson production in hadronic collisions, when suitably binned, provides unambiguous tests of whether the Pomeron couples like a conserved vector-current to protons. We discuss the possibility of discriminating between $q\bar{q}$ and glueball production in such processes. Our predictions apply also to meson production in tagged two-photon events at electron–positron colliders and to vector-meson production in ep collisions at HERA.

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1 Introduction

The production of mesons in the central region of proton–proton collisions (pp $\to$ pp$M$) via a gluonic Pomeron has traditionally been regarded as a potential source of glueballs [1]. However, well-established quark–antiquark ($q\bar{q}$) mesons are also known to be produced and this has led to searches for a selection mechanism that could help to distinguish among such states. As a result it has been discovered [2] that the pattern of resonances produced in the central region of double tagged pp $\to$ pp$M$ depends on the vector *difference* $\vec{k}_\perp = \vec{q}_{1\perp} - \vec{q}_{2\perp}$ of the transverse momentum recoils $\vec{q}_{i\perp}$ of the final protons (even at fixed four-momentum transfers $t_i = -q_i^2$). When this quantity $k_{\rm T} = |\vec{k}_\perp|$ is small ($\leq \mathcal{O}(\Lambda_{QCD})$) all well-established $q\bar{q}$ states were observed to be suppressed [3] while the surviving resonances included enigmatic states such as $f_0(1500)$, $f_J(1710)$ and $f_2(1910)$ that have variously been suggested to be glueballs or to reside on the (gluonic) Pomeron trajectory. At large $k_{\rm T}$, by contrast, $q\bar{q}$ states are prominent.

However, these $k_{\rm T}$ dependences for at least 0^- and 1^+ production have been shown to arise if the Pomeron (or perhaps a hard gluonic component that produces M by gg fusion) transforms as a conserved vector current [4]. In order to help determine the extent to which the double tagged reaction pp $\to$ pp$M$ depends on a vector production or the dynamical structure of the meson M (of spin J and parity P), we develop the earlier analysis to all $J \leq 3$.

While the $k_{\rm T}$ phenomenon has turned out to be a sharp experimental signature, we shall propose here that the azimuthal ϕ dependence (between the two proton scattering-planes in the pp c.m.s.) provides a rather direct probe of dynamics. In particular, observation of non-trivial ϕ dependences requires the presence of non-zero helicity transfer by the diffractive agent (Pomeron, gluon, ...) [5] and so the Pomeron cannot simply transform as having vacuum quantum numbers: a spin greater than zero is needed. We analyse here the simplest case, where the process is driven by the fusion of two spin-1 currents. Imposing current conservation it immediately applies to $e^+e^- \to e^+e^- M$ and, empirically, already exhibits features seen in pp $\to$ pp$M$. We find that current non-conserving and/or scalar contributions are needed to accommodate the data.

At extreme energies where non-diffractive contributions are negligible, we show the following properties for meson production in the central region.

1. The ϕ dependence of 0^- production provides a clear test for the presence of a significant vector component of the production Pomeron, independent of the t dependence. Preliminary data on η and η' production confirm this.

2. The production of 1^+ mesons reinforces this: The conserved vector-current (CVC) hypothesis implies (i) the cross section will tend to zero as $k_{\rm T} \to 0$, and (ii) 1^+ mesons are produced dominantly in the helicity-one state. Both features are prominent in the data.

3. The 0^+ cross section survives at small $k_{\rm T}$ for the CVC hypotheses. Moreover, for $q_{i\perp} \ll M$, we must either observe a $\cos^2\phi$ distribution or a small (relative to

0^-) cross section. However, at least one 0^+ state (the $f_0(2000)$) appears to be suppressed at small k_{T}. Unlike 0^- or 1^+ production, the production of 0^+ will be particularly sensitive to a scalar and/or non-conserved vector component to the Pomeron. In particular the vanishing of $f_0(1500)$ as $\phi \to 180^o$ like $\sin^4(\phi/2)$, would be natural if longitudinal and transverse helicity amplitudes have similar strengths but opposite phase as may be possible in some simple glueball models.
4. The 2^+ production depends on the dynamics of the meson as well as the helicity structure of the Pomeron. In the non-relativistic $\mathrm{q\bar{q}}$ model (a particular realization of the CVC hypothesis), we predict at small $q_{i\perp} \ll M$ a 2^+ cross section that is (i) basically flat in $\cos\phi$, (ii) finite for $k_{\mathrm{T}} \to 0$, and (iii) dominated by the helicity-two part. For the CVC hypothesis a suppression at small k_{T} is obtained only for peculiar relations between the helicity amplitudes. Hence again, data show that the CVC Pomeron is not the full story. In particular, 2^+ states at 2 GeV are seen to have a different ϕ dependence than the established $\mathrm{q\bar{q}}$ 2^+ states.

Our analysis can also be applied to $\mathrm{ep} \to \mathrm{ep}\, M$ where M is a vector meson. As $t \to 0$ we find that the longitudinal polarization of the meson grows initially as Q^2/M^2 relative to the transverse, with a characteristic ϕ dependence.

Readers interested in the results may proceed directly to section 3. Their detailed derivation is summarized in section 2.

2 Derivation of the results

Consider the central production of a J^{P+} meson M in the high-energy scattering of two fermions with momenta p_1 and p_2, respectively,

$$f_1(p_1) + f_2(p_2) \to f_1'(p_1') + f_2'(p_2') + M \ , \tag{1}$$

proceeding through the fusion of two conserved spin-1 vector currents V_1 and V_2:

$$V_1(q_1, \lambda_1) + V_2(q_2, \lambda_2) \to M(J, J_z) \ . \tag{2}$$

Here $\lambda_i = \pm 1$, L are the current helicities[1] in the meson rest frame with current one defining the z axis. In the case of electron–positron collisions, V_i in (2) is a photon, while for central production in proton–proton collisions, V_i could be a Pomeron, a (colour-less) multi-gluon state, or in some models, even a single gluon (accompanied by Coulomb gluon(s) to ensure colour-conservation). For our purpose here what matters is the assumed spin-1 nature of the production field(s) and their conservation. We shall comment upon the consequences of current non-conservation at the end of the next section.

[1] Our longitudinal-helicity polarization vector $\epsilon_\mu(L)$ is orthogonal to the momentum vector as are the two transverse polarization vectors. Hence in the meson rest frame $\epsilon_\mu(L)$ has both a 0 and a 3 component. Consequently, our scalar polarization vector is proportional to the momentum vector.

In order to investigate the helicity structure of the diffractive agent it proves useful to examine the dependence of the cross section on[2] $\cos n\tilde{\phi}$, where $\tilde{\phi}$ is the azimuthal separation between the two proton scattering planes of (1) in the current-current c.m.s. An experimental analysis is complicated by two facts. First, what is measured is not $\tilde{\phi}$ but the azimuthal angle ϕ in the proton–proton rest frame. Second, experimental cuts and/or an inconvenient choice of kinematical variables might spoil the $\tilde{\phi}$ dependence predicted by theory.

This is easily understood when one recalls that the phase space for (1), $\sim (\mathrm{d}^3p'_1/E'_1)\,(\mathrm{d}^3p'_2/E'_2)$, depends on only four non-trivial variables if the meson is either stable or has a width much smaller than its mass since one relation is provided by $W \equiv \sqrt{(q_1+q_2)^2} = M$ (the meson mass). These four variables are often chosen as four invariants, for example, $Q_i = \sqrt{-q_i^2}$ and (suitably defined) fractional current energies x_i, or as the scattered protons energies and polar angles (or transverse momenta). For whatever choice, the expression of (the fixed variable) W in terms of these variables explicitly involves the angle $\tilde{\phi}$ (or ϕ), which introduces additional "spurious" azimuthal dependences. Moreover, the relation between $\tilde{\phi}$ and the measurable ϕ is rather complicated.

However, as we shall detail below, in the kinematic regime of experimental interest, we have, to good approximation, $\tilde{\phi} \approx \phi$. Also for the WA102 experiment, we estimate the effect of the extra kinematic factors to have no significant impact on the effects discussed here.

In the approximation of single-particle (single-trajectory) exchange (one at each vertex) the cross section for (1) factors into the product of three terms, namely two density matrices and the amplitude for (2). Consider the (unnormalized) density matrix for the emission from particle 1. For a conserved vector-current its general form is

$$\rho_1^{\mu\nu} = -\left(g^{\mu\nu} - \frac{q_1^\mu\, q_1^\nu}{q_1^2}\right) C_1(q_1^2) - \frac{(2p_1-q_1)^\mu\,(2p_1-q_1)^\nu}{q_1^2}\, D_1(q_1^2)\ . \qquad (3)$$

Here C_1 and D_1 are form factors associated with the non-pointlike nature of particle 1 (for a lepton, $C_\mathrm{e}(q^2) = 1 = D_\mathrm{e}(q^2)$, while for a proton $C_\mathrm{p}(q^2) = G_\mathrm{M}^2(q^2)$, $D_\mathrm{p}(q^2) = (4m_\mathrm{p}^2 G_\mathrm{E}^2(q^2) - q^2 G_\mathrm{M}^2(q^2))/(4m_\mathrm{p}^2 - q^2)$, where G_E and G_M are the proton electromagnetic form factors). A factor $1/(-2q_1^2)$ in ρ_1 is introduced for convenience[3] since current conservation guarantees $(2p_1-q_1)^\mu(2p_1-q_1)^\alpha\, M^{\star\alpha\beta} M^{\mu\nu} \propto q_1^2$.

In the following we shall be working in the current–current helicity basis. The density-matrix elements in the helicity basis are defined with the help of the polarization vectors $\epsilon_1^\mu(\lambda_1)$ of the (space-like) current one as [6]

$$\rho_1^{\lambda_1,\lambda'_1} = (-1)^{\lambda_1+\lambda'_1}\, \epsilon_1^\mu(\lambda_1)\, \rho_1^{\mu\nu}\, \epsilon_1^\nu(\lambda'_1)\ , \qquad (4)$$

[2] P and T invariance forbid $\sin n\tilde{\phi}$ contributions.

[3] With this choice, the matrix elements of the first term are simply $\pm C_1$ (or zero), see (6) and (12).

where $\lambda_1^{(\prime)}$ label the helicity of the current one, $\lambda_1^{(\prime)} = \pm 1, L$. Owing to the hermiticity relations of the density matrix and the polarization vectors

$$\begin{aligned} \rho_1^{\mu\nu\star} &= \rho_1^{\nu\mu} \\ \epsilon_1^{\alpha\star}(\pm 1) &= -\epsilon_1^{\alpha}(\mp 1)\ , \qquad \epsilon_1^{\alpha\star}(L) = -\epsilon_1^{\alpha}(L)\ , \end{aligned} \tag{5}$$

the helicity-density matrix is determined by four real parameters, for example, ρ_1^{++}, ρ_1^{LL}, $|\rho_1^{+L}|$, and $|\rho_1^{+-}|$. The phases of the latter two matrix elements are $\exp(i\tilde{\phi}_1)$ and $\exp(2i\tilde{\phi}_1)$, respectively, where $\tilde{\phi}_1$ is the azimuthal angle of p_1 in the current–current c.m.s. (With the analogous definition of $\tilde{\phi}_2$ we have $\tilde{\phi} = \tilde{\phi}_1 + \tilde{\phi}_2$.)

The expressions of $|\rho_1^{ik}|$ in terms of invariants and the form factors C_1 and D_1 can be derived from the formulas in [6]

$$\begin{aligned} \rho_1^{++} &= C_1 + \frac{1}{2} D_1 \left[\frac{(u_2 - \nu)^2}{X} - 1 + \frac{4\, m_1^2}{q_1^2} \right] \\ \rho_1^{LL} &= -C_1 + D_1 \frac{(u_2 - \nu)^2}{X} \\ |\rho_1^{+-}| &= \rho_1^{++} - C_1 \\ |\rho_1^{+L}| &= \sqrt{|\rho_1^{+-}|\,(\rho_1^{LL} + C_1)}\ . \end{aligned} \tag{6}$$

Here we have introduced $u_2 = 2 p_1 \cdot q_2$, $\nu = q_1 \cdot q_2 = (W^2 - q_1^2 - q_2^2)/2$, $W^2 = (q_1 + q_2)^2$, and $X = \nu^2 - q_1^2\, q_2^2$.

In this work we are interested in the dominant (and experimentally accessible) region of phase space $Q_i \equiv \sqrt{-q_i^2} \ll W$. Then (and only then [7]) the density matrix $\rho_1^{\mu\nu}$ depends on only variables of current-one, namely its fractional momentum $x_1 = p_2 \cdot q_1 / p_2 \cdot p_1 = u_1/(s - 2m_1^2)$ and its virtuality Q_1. Moreover, we can use

$$\begin{aligned} Q_i &\simeq q_{i\perp} \\ \tilde{\phi} &= \frac{\vec{q}_{1\perp} \cdot \vec{q}_{2\perp}}{q_{1\perp}\, q_{2\perp}} \simeq \phi = \frac{\vec{p}^{\,\prime}_{1\perp} \cdot \vec{p}^{\,\prime}_{2\perp}}{p'_{1\perp}\, p'_{2\perp}}\ , \end{aligned} \tag{7}$$

where $\vec{q}_{i\perp}$ ($\vec{p}^{\,\prime}_{i\perp}$) is the transverse momentum of current i (scattered proton i) in the current-current (proton–proton) c.m. system. In addition, the dependence of $W = M$ on the azimuthal angle $\tilde{\phi} = \tilde{\phi}_1 + \tilde{\phi}_2$ disappears, and we simply have $W^2 = x_1 x_2 s$ ($x_2 = u_2/(s - 2m_2^2)$). Since m_1 and m_2 are much smaller than the c.m. energy $\sqrt{s}$ we obtain

$$\begin{aligned} 2\,\rho_1^{++} &= 2\,C_1 + (1 - \delta_1)\,\hat{\rho}_1 \\ \rho_1^{LL} &= D_1 - C_1 + \hat{\rho}_1 \\ 2\,|\rho_1^{+-}| &= (1 - \delta_1)\,\hat{\rho}_1 \\ \sqrt{2}|\rho_1^{+L}| &= \sqrt{(1 - \delta_1)\,\hat{\rho}_1\,(D_1 + \hat{\rho}_1)}\ , \end{aligned} \tag{8}$$

where we have introduced

$$\hat{\rho}_1 = \frac{4}{x_1^2}\,(1-x_1)\;D_1\;, \qquad \delta_1 = Q_{1\,\mathrm{min}}^2/Q_1^2\;. \tag{9}$$

For the production (1) of mesons at fixed-target experiments (and even more so at electron–positron colliders) the meson mass is much smaller than the c.m. energy. This implies that $x_i \ll 1$ and thus

$$\frac{2}{1-\delta_1}\,\rho_1^{++} \simeq \frac{2}{1-\delta_1}\,|\rho_1^{+-}| \simeq \sqrt{\frac{2}{1-\delta_1}}\,|\rho_1^{+L}| \simeq \rho_1^{LL} \simeq \hat{\rho}_1\;. \tag{10}$$

Relations analogous to (6)–(10) hold also for the density matrix of current two, $\rho_2^{\lambda_2,\lambda_2'}$.

Before continuing we have to make sure that (10) is not spoiled by the behaviour of the form factors, i.e. we have to make sure that $\hat{\rho}_1 \gg C_1$. This is certainly true if $C_1 \simeq D_1$ for all Q_1^2. To investigate this a bit further we assume that Pomerons (e.g. Pomeron one) couple to fermions like the current

$$J_\mu = \bar{u}(p_1')\,\left\{F_1(q_1^2)\,\gamma_\mu + \frac{\kappa}{2m}\,F_2(q_1^2)\,i\sigma_{\mu\alpha}\,q^\alpha\right\}\,u(p_1) \tag{11}$$

Then we can actually calculate the density matrix (3) defined by

$$\rho_1^{\mu\nu} = \frac{-1}{2\,q_1^2}\sum_{\mathrm{spins}} J_\mu\,J_\nu^\star\;. \tag{12}$$

Noting the minus sign in

$$(\bar{u}(p_1')\,i\sigma_{\mu\alpha}\,q^\alpha\,u(p_1))^\star = -\bar{u}(p_1)\,i\sigma_{\mu\alpha}\,q^\alpha\,u(p_1')\;,$$

we obtain the form (3) with

$$\begin{aligned} C_1 &= (F_1+\kappa\,F_2)^2 \equiv G_{\mathrm{M}}^2 \\ D_1 &= F_1^2 - \frac{q_1^2}{4m^2}\,(\kappa\,F_2)^2 \equiv \frac{4m^2\,G_{\mathrm{E}}^2 - q_1^2\,G_{\mathrm{M}}^2}{4m^2-q_1^2}\;. \end{aligned} \tag{13}$$

Note that a pure γ_μ coupling gives $C_1 = D_1 = F_1^2$. Hence for two-photon production at $\mathrm{e}^+\mathrm{e}^-$ colliders ($F_1 = 1$, $F_2 = 0$) our assumptions are well satisfied: even at CLEO energies the typical x_i values are small enough ($x_i \sim 0.1$) to ensure $1/x_i^2 \gg 1$ and, in turn, $\hat{\rho}_1 \gg C_1$. Moreover, the tagging setup of the scattered electrons assures that $\delta_i \ll 1$.

The situation may be different in fixed-target proton–proton collisions. First, at WA102 energies ($12.8 < \sqrt{s}/\mathrm{GeV} < 28$) the experimentally accessible x_i values range between about 10^{-3} and 0.2 guaranteeing thus $1/x_i^2 \gg 1$. The minimum x_i values result in minimum virtualities of $Q_{i\,\mathrm{min}}^2 \approx 10^{-4}\,\mathrm{GeV}^2$. Hence if we assume that measurements are done in a range, say $10^{-3} < Q_i^2/\mathrm{GeV}^2 < 0.5$ (statistics limits larger values) then still $\delta_i \ll 1$. This holds certainly for the recoil proton since it can only be detected for Q^2 larger than about $0.05\,\mathrm{GeV}^2$. The scattered proton can, however, be measured down to very low Q^2. For completeness, we shall keep the $(1-\delta_i)$ terms in the following.

Central production in proton–proton collisions may differ in another aspect from the e^+e^- case: unlike the photon the Pomeron might have a dominant $\sigma_{\mu\nu}$-type coupling. The requirement for (10) to hold, namely $\hat{\rho}_1 \gg C_1$, yields for zero F_1 the condition $Q_1^2 \gg Q_{1\,\mathrm{min}}^2$. Hence as long as very low Q^2 values of the scattered proton are excluded, (10) continues to hold. There is one difference, however: if F_2 dominates then the typical t ($t = q^2 = -Q^2$) distribution $\propto \exp(-bt)$ (with $b \sim 6/\mathrm{GeV}^2$) is modified by an extra factor $(-t)$.

Let us now continue with the current–current–meson vertex. If (2) proceeds through the fusion of two conserved vector-currents, then conservation of P and T as well as total helicity conservation for forward scattering, implies that the cross section for $f_1 + f_2 \to f_1' + f_2' + X$, for arbitrary final state X, depends on eight independent helicity structure functions, $W(\lambda_1, \lambda_2; \lambda_1', \lambda_2')$ out of which only six can be measured with unpolarized initial-state fermions:

$$\begin{aligned} \mathrm{d}\sigma \sim\ & 2\,\rho_1^{++}\,\rho_2^{++}\,\{W(++,++) + W(+-,+-)\} \\ & + 2\,\rho_1^{++}\,\rho_2^{LL}\,W(+L,+L) \\ & + 2\,\rho_1^{LL}\,\rho_2^{++}\,W(L+,L+) \\ & + \rho_1^{LL}\,\rho_2^{LL}\,W(LL,LL) \\ & + 2\,\left|\rho_1^{+-}\,\rho_2^{+-}\right|\,W(++,--)\,\cos 2\tilde{\phi} \\ & - 4\,\left|\rho_1^{+L}\,\rho_2^{+L}\right|\,\{W(++,LL) + W(L+,-L)\}\,\cos\tilde{\phi}\ . \end{aligned} \tag{14}$$

Note that $W(\lambda_1, \lambda_2; \lambda_1', \lambda_2') \neq 0$ only if $\lambda_1 - \lambda_2 = J_z = \lambda_1' - \lambda_2'$. Both the structure functions W and the invariant amplitudes A defined below in (15) are functions of the invariants W, Q_1^2, Q_2^2 only[4].

For the present case, (1), where X is a single particle, the number of independent parameters in (14) can be reduced further. First observe that if $A(\lambda_1, \lambda_2)$ denotes the ($V_1 V_2$ c.m.s.) helicity amplitude for (2), then we have

$$W(\lambda_1, \lambda_2; \lambda_1', \lambda_2') = A(\lambda_1, \lambda_2)\,A^\star(\lambda_1', \lambda_2')\,\delta\left(W^2 - M^2\right)\ , \tag{15}$$

where $W^2 = (q_1 + q_2)^2$ and M denotes the meson mass. Second, if η_i denotes the naturality[5] of current V_i and η_M that of the meson M, then

$$A(-\lambda_1, -\lambda_2) = \eta\,A(\lambda_1, \lambda_2) \quad , \quad \eta \equiv \eta_1\,\eta_2\,\eta_M\ , \tag{16}$$

[4] If instead one chose to replace one of these variables by ϕ then different ϕ dependences could emerge, see, for example, (5.14) in [6] or (13) in [8].

[5] A boson is said to have naturality $+1$ if $P = (-1)^J$ and -1 if $P = (-1)^{J-1}$.

and there are five independent helicity amplitudes $A(\lambda_1,\lambda_2)$. Finally, owing to the T-invariance relation

$$W(\lambda_1,\lambda_2;\lambda_1',\lambda_2') = W(\lambda_1',\lambda_2';\lambda_1,\lambda_2)\ , \tag{17}$$

which implies

$$A(\lambda_1,\lambda_2)\,A^\star(\lambda_1',\lambda_2') = \mathrm{csgn}\,A(\lambda_1,\lambda_2)\,\mathrm{csgn}\,A(\lambda_1',\lambda_2')\ |A(\lambda_1',\lambda_2')|\ |A(\lambda_1,\lambda_2)|\ , \tag{18}$$

we are left with five real parameters. Here

$$\mathrm{csgn}\,z = \begin{cases} +1 & \mathrm{Re}\,z > 0\ \text{ or } (\mathrm{Re}\,z = 0 \text{ and } \mathrm{Im}\,z > 0) \\ -1 & \mathrm{Re}\,z < 0\ \text{ or } (\mathrm{Re}\,z = 0 \text{ and } \mathrm{Im}\,z < 0)\ . \end{cases} \tag{19}$$

Defining $A_{\lambda_1\lambda_2} = |A(\lambda_1,\lambda_2)|$ and

$$\begin{aligned} \xi_1 &= \mathrm{csgn}\,A(++)\,\mathrm{csgn}\,A(LL) \\ \xi_2 &= \mathrm{csgn}\,A(+L)\,\mathrm{csgn}\,A(L+)\ , \end{aligned} \tag{20}$$

we find

$$\begin{aligned} \mathrm{d}\sigma \ \sim\ & 2\,\rho_1^{++}\,\rho_2^{++}\,A_{+-}^2 \\ & + 2\,\rho_1^{++}\,\rho_2^{LL}\,A_{+L}^2 + 2\,\rho_1^{LL}\,\rho_2^{++}\,A_{L+}^2 - 4\,\eta\,\left|\rho_1^{+L}\,\rho_2^{+L}\right|\,\xi_2\,A_{+L}\,A_{L+}\,\cos\tilde\phi \\ & + \rho_1^{LL}\,\rho_2^{LL}\,A_{LL}^2 - 4\,\left|\rho_1^{+L}\,\rho_2^{+L}\right|\,\xi_1\,A_{++}\,A_{LL}\,\cos\tilde\phi \\ & + \left\{2\,\rho_1^{++}\,\rho_2^{++} + 2\,\eta\,\left|\rho_1^{+-}\,\rho_2^{+-}\right|\,\cos 2\,\tilde\phi\right\}\,A_{++}^2\ . \end{aligned} \tag{21}$$

For the kinematic regime of interest, $x_i \ll 1$ and $Q_1 \ll M$, $\tilde\phi \approx \phi$, (7), and (10) allows us to approximate

$$\begin{aligned} \rho_1^{LL}\,\rho_2^{LL} &\approx \frac{4\,\rho_1^{++}\,\rho_2^{++}}{(1-\delta_1)\,(1-\delta_2)} \approx \frac{2\,\rho_1^{++}\,\rho_2^{LL}}{1-\delta_1} \approx \frac{2\,\rho_1^{LL}\,\rho_2^{++}}{1-\delta_2} \\ &\approx \frac{2\,\left|\rho_1^{+L}\,\rho_2^{+L}\right|}{\sqrt{(1-\delta_1)\,(1-\delta_2)}} \approx \frac{4\,\left|\rho_1^{+-}\,\rho_2^{+-}\right|}{(1-\delta_1)\,(1-\delta_2)}\ . \end{aligned} \tag{22}$$

If we decompose the cross section into components (subscript i on Σ_i) that correspond to $|J_z| = 2$, 1, and 0, then we obtain

$$\begin{aligned} \mathrm{d}\sigma &\sim \Sigma_2 + \Sigma_1 + \Sigma_0 \\ \Sigma_2 &= \frac{1}{2}\,(1-\delta_1)\,(1-\delta_2)\,A_{+-}^2 \\ \Sigma_1 &= (1-\delta_1)\,A_{+L}^2 + (1-\delta_2)\,A_{L+}^2 \\ &\quad - 2\,\eta\,\xi_2\,\sqrt{(1-\delta_1)\,(1-\delta_2)}\,A_{+L}\,A_{L+}\,\cos\phi \\ \Sigma_0 &= A_{LL}^2 - 2\,\xi_1\,\sqrt{(1-\delta_1)\,(1-\delta_2)}\,A_{++}\,A_{LL}\,\cos\phi \\ &\quad + (1-\delta_1)\,(1-\delta_2)\,(1+\eta\,\cos 2\,\phi)\,\frac{1}{2}A_{++}^2\ . \end{aligned} \tag{23}$$

J^P	0^-	0^+	1^-	1^+	2^-	2^+	3^-	3^+
η	$-$	$+$	$+$	$-$	$-$	$+$	$+$	$-$
A_{LL}	0	δ	$D\,\delta$	0	0	δ	$D\,\delta$	0
A_{++}	1	1	D	D	1	1	D	D
A_{+-}	0	0	0	0	D	1	D	1
κ	1	0	1	0	1	0	1	0

Table 1: Model-independent features of helicity amplitudes up to $J^P = 3^+$; 0: amplitude is identical to zero; D: amplitude is proportional to $D = (Q_1^2 - Q_2^2)/M^2$; δ: amplitude is proportional to $\delta = Q_1\,Q_2/M^2$ for $Q_i \ll M$; 1: amplitude is of order one, in general. Also given are the values of η, (16), and κ, (29) (for the case $\eta_1\,\eta_2 = +1$).

Introducing

$$r = \frac{A_{LL}}{A_{++}} \ , \tag{24}$$

and making use of $(1-\eta)\,r = 0$, we can rewrite the $J_z = 0$ part in (23) as

$$\begin{aligned} \Sigma_0 &= A_{++}^2 \left\{ \delta_{\eta,1} \left(r - \xi_1 \sqrt{(1-\delta_1)\,(1-\delta_2)}\,\cos\phi \right)^2 \right. \\ &\qquad \left. + \, \delta_{\eta,-1}\,(1-\delta_1)\,(1-\delta_2)\,\frac{1}{2}\,(1-\cos 2\,\phi) \right\} \ . \end{aligned} \tag{25}$$

Which of the two terms in (25) contributes depends on the naturality factor η, see table 1.

So far we have not yet made use of Bose symmetry, which states

$$A(\lambda_1,\lambda_2)(Q_1,Q_2) = (-1)^J\,A(\lambda_2,\lambda_1)(Q_2,Q_1) \ , \tag{26}$$

where $Q_i = \sqrt{-q_i^2}$ is the virtuality of boson i. It implies that (in the CVC hypothesis) the amplitudes A_{++} and A_{LL} must be proportional to

$$D = \frac{Q_1^2 - Q_2^2}{M^2} \tag{27}$$

for odd-integer J. When combined with parity, (16), the amplitude $A_{+-} \propto D$ for some J^P, see table 1.

Bose symmetry has one more consequence, namely that both amplitudes A_{L+} and A_{+L} in (23) can be replaced by only one of them, say A_{+L}. Moreover, the sign in (18) is then fixed in a model-independent way. We can rewrite the $|J_z| = 1$ part of the cross section as

$$\begin{aligned} \Sigma_1 &= (1-\delta_1)\,A_{+L}^2(Q_1,Q_2) + (1-\delta_2)\,A_{+L}^2(Q_2,Q_1) \\ &\quad -2\,(1-2\,\kappa)\sqrt{(1-\delta_1)\,(1-\delta_2)}\,A_{+L}(Q_1,Q_2)\,A_{+L}(Q_2,Q_1)\,\cos\phi \ , \end{aligned} \tag{28}$$

where we have introduced the variable

$$\kappa = \frac{1 - \eta\,(-1)^J}{2}\ , \tag{29}$$

whose values, one or zero, are given in table 1 for states up to $J^P = 3^+$.

We can exploit one more constraint, namely current conservation, which requires

$$\begin{aligned} A_{\pm 1,L} &\propto \quad Q_2/M \qquad &&\text{for} \quad Q_2 \ll M \\ A_{L,\pm 1} &\propto \quad Q_1/M \qquad &&\text{for} \quad Q_1 \ll M \\ A_{L,L} &\propto \ Q_1\,Q_2/M^2 \qquad &&\text{for} \quad Q_i \ll M\ . \end{aligned} \tag{30}$$

Then (7) implies that

$$\begin{aligned} A_{+L} &\simeq a_{+L}\,\frac{q_{2\perp}}{M} \\ A_{LL} &\simeq a_{LL}\,\delta\ , \qquad \delta \equiv \frac{Q_1\,Q_2}{M^2} \simeq \frac{q_{1\perp}\,q_{2\perp}}{M^2} \\ D &\simeq \frac{q_{1\perp}^2 - q_{2\perp}^2}{M^2}\ , \end{aligned} \tag{31}$$

where a_{ij} are coefficients of order one. Hence Σ_1 in (28) behaves as

$$\Sigma_1 = \left\{ \begin{array}{ll} a_{+L}^2\,p_{\rm T}^2/M^2\,, & \text{for } \kappa = 1\ , \\ a_{+L}^2\,k_{\rm T}^2/M^2\,, & \text{for } \kappa = 0\ , \end{array} \right. \tag{32}$$

where

$$\begin{aligned} p_{\rm T}^2 &= \left(\sqrt{1-\delta_2}\,\vec{q}_{1\perp} + \sqrt{1-\delta_1}\,\vec{q}_{2\perp}\right)^2 \\ k_{\rm T}^2 &= \left(\sqrt{1-\delta_2}\,\vec{q}_{1\perp} - \sqrt{1-\delta_1}\,\vec{q}_{2\perp}\right)^2\ . \end{aligned} \tag{33}$$

Note that $k_{\rm T} \to 0$ implies $\phi \to 0$ *and* $q_{2\perp} \to q_{1\perp}$. However, the opposite is not true: $\phi \to 0$ does not in general imply $k_{\rm T} \to 0$.

3 Results

The above analysis enables some immediate conclusions to be drawn according to the J^{PC} of the meson.

(i) J^P=0$^-$

Only $J_z = 0$ contributes and, with $\eta = -1$ in (25)

$$\frac{d\sigma}{d\phi} \propto A_{++}^2\,(1-\delta_1)\,(1-\delta_2)\,\frac{1}{2}\,(1 - \cos 2\,\phi) = A_{++}^2\,(1-\delta_1)\,(1-\delta_2)\,\sin^2\phi\ . \tag{34}$$

This follows independent of the dynamical internal structure of the 0^{-+} meson, and is simply a consequence of parity. Since $\phi \to 0$ as $k_{\rm T} \to 0$ we recover the result of [4, 9, 10] who noted that the production of 0^{-+} by (conserved) vector currents would vanish as $k_{\rm T} \to 0$. Our result above provides a clear test for the vector nature of the production Pomeron (component) by the explicit prediction for the ϕ dependence, independent of the t-dependence.

Preliminary indications are that the production of η and η' in pp $\to$ ppη (η') is compatible with such a ϕ dependence [11].

(iii) J^P=1$^+$

Since $|J_z| \leq 1$ the azimuthal distribution is given by the sum of (25) (with $\eta = -1$) and (28) (with $\kappa = 0$). Since Bose symmetry yields $A_{++} = a_{++}\, D$ we find with the help of table 1 and (31) in the region of small Q_i

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\phi} \sim a_{+L}^2\, \frac{k_{\rm T}^2}{M^2} + a_{++}^2\, (1-\delta_1)\,(1-\delta_2)\, \sin^2\phi\, \frac{(q_{1\perp}^2 - q_{2\perp}^2)^2}{M^4}\,. \qquad (35)$$

From this we can draw conclusions, which are independent of the internal structure of the 1^{++} meson and thus hold for both e$^+$e$^-$ collisions and diffractive proton–proton collisions mediated by a vector Pomeron. First, the cross section will tend to zero as $k_{\rm T} \to 0$. And second, 1^+ mesons are produced dominantly in the helicity-one state. Both of these phenomena are seen in the central production of 1^{++} mesons in pp collisions which further supports the importance of the vector component of the effective Pomeron.

The tendency for large $k_{\rm T}$ to correlate with large ϕ may cause the apparent $\mathrm{d}\sigma/\mathrm{d}\phi$ to rise as $\phi \to 180^o$. The ϕ distributions should be binned in $k_{\rm T}$ to extract the full implications of (35).

(ii) J^P=0$^+$

In this case the ϕ dependence depends on the internal structure of the meson and dynamics, specifically via the magnitude of $A_{LL}/A_{++} \equiv r$

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\phi} = A_{++}^2\, \left(\sqrt{(1-\delta_1)\,(1-\delta_2)}\,\xi_1\, \cos\phi - r\right)^2\,. \qquad (36)$$

At small Q_i, $Q_i \ll M$ in e$^+$e$^-$ collisions, we have $r \simeq c\,\delta \simeq c\, q_{1\perp}\, q_{2\perp}/M^2$ with $c = a_{LL}/A_{++} = O(1)$, in general.

For the particular case of two photons coupling to a non-relativistic quark-antiquark one has [12, 13] $\xi_1 = +1$ and $c = 4/3$ since

$$r = \frac{Q_1\, Q_2\, M^2}{\nu^2 + \nu\, M^2 - Q_1^2\, Q_2^2} \approx \frac{4}{3}\, \frac{q_{1\perp}\, q_{2\perp}}{M^2} \qquad (37)$$

at $q_{i\perp} \ll M$. Hence for tagged two-photon events in e$^+$e$^-$ collisions we predict a cross section that survives the $k_{\rm T} \to 0$ limit and the ϕ distribution (36), which for $q_{i\perp} \ll M$ is a pure $\cos^2\phi$ distribution.

This will also hold true for $q\bar{q}$ *and* glueball production in pp collisions *if* the Pomeron is a *conserved* vector current. So far we have taken the simplest assumption needed for a nontrivial ϕ distribution, namely CVC. This is immediately relevant to e^+e^- but encouragingly shows consistency with pp. The 0^- is a direct test with its $\sin^2\phi$ distribution which is verified for η, η' in WA102. For 1^+ the $k_T \to 0$ vanishing and the helicity-1 dominance are verified. The 0^+, 2^+ data clearly go beyond this.

The non-trivial ϕ dependence required $J_{\rm Pomeron} > 0$ to be present but leaves open the question of whether there is a spin-0 component in addition to the CVC and/or whether there is a non-conserved vector current. Note that the 0^- production is not sensitive to any 0^+ component in the Pomeron. The simplest manifestation of a scalar component or a non-conserved vector piece, is to allow R to be larger than its CVC suppression $O(\sqrt{t_1 t_2}/M^2)$. The 0^+, 2^+ data are consistent with this. The $f_0(1500)$ production, in particular, is well described if R is negative with $|R| \sim O(1)$, in which case its ϕ distribution is $\sim \sin^4(\phi/2)$. This sign and magnitude are natural for the production of a gluonic system if the dynamics for M_{LL}/M_{++} is driven by the Clebsch–Gordon coefficients $\langle 10,10|00\rangle/\langle 11,1-1|00\rangle = -1$. We leave the discussion of the phenomenology and specific models to a later publication.

(iii) J^P=2$^+$

The azimuthal distribution is given by the sum of Σ_2, (23), Σ_1, (28) with $\kappa = 0$, and Σ_0, (25) with $\eta = +1$. Using the small-Q_i approximation for Σ_1 we have

$$d\sigma \sim \frac{1}{2}\,A_{+-}^2\,(1-\delta_1)\,(1-\delta_2) + a_{+L}^2\,\frac{k_T^2}{M^2} + \left(r - \xi_1\,\sqrt{(1-\delta_1)\,(1-\delta_2)}\,\cos\phi\right)^2\,A_{++}^2\,. \tag{38}$$

As we can see the $|J_z| = 1$ part is suppressed as is A_{LL} (recall $r \sim \delta$ at small Q_i). However, in general, the other two amplitudes are of order one, i.e. $A_{+-} \sim A_{++} \sim 1$.

In the non-relativistic quark model, $A_{++} \simeq (Q_1^2 + Q_2^2)/M^2$ at small Q_i [14, 12] and is thus very much suppressed relative to A_{+-}, which is $O(1)$. Hence in e^+e^- collisions at small $q_{i\perp} \ll M$ we predict a 2^+ cross section that is (i) basically flat in $\cos\phi$, (ii) finite for $k_T \to 0$, and (iii) dominated by the helicity-two part. We necessarily obtain the same behaviour, namely flat ϕ distribution and $k_T \to 0$ survival, in diffractive pp collisions mediated by a conserved vector Pomeron, provided the helicity-two component is the dominant one.

If the Pomeron–$q\bar{q}$ coupling were dominantly "magnetic" (flipping the spins of the produced $q\bar{q}$ pair but leaving them in an $L_z = 0$ state) the helicity-two amplitude A_{+-} would be suppressed. In this case the helicity-one amplitude would also be suppressed as $k_T \to 0$ and the helicity-zero amplitude would dominate with a characteristic ϕ dependence (unless $A_{++} = 0$). Moreover, the 2^+ cross section continues to survive the $k_T \to 0$ limit since r is small for CVC.

Again we conclude that, as for 0^+ production, a non-conserved vector piece (or a large scalar component) is needed to accommodate for the observed small-$k_{\rm T}$ suppression of $\rm f_2(1270)$ and $\rm f_2'(1520)$. In the scenario discussed above this follows if $\xi_1\, r \sim O(1)$. We point out that these predictions assume Pomeron–Pomeron or gluon–gluon fusion and hence do not apply to $\rm f_2$ production if the latter has a substantial contribution from $\rm f_2$ exchange (i.e. from $\rm f_2$+ Pomeron $\to \rm f_2$. A detailed comparison of $\rm f_2\ s\bar{s}(1525)$ (for which this contribution is suppressed) and $\rm f_2\ u\bar{u}(1270)$ (where Pomeron–$\rm f_2$ is possible) could help settle this.

(iii) J^P=2$^-$

Here we find with the help of table 1 and (31)

$$
{\rm d}\sigma \sim (1-\delta_1)\,(1-\delta_2)\left\{\frac{1}{2}\,a_{+-}^2\,\frac{\left(q_{1\perp}^2-q_{2\perp}^2\right)^2}{M^4}+\sin^2\phi\,A_{++}^2\right\}+a_{+L}^2\,\frac{p_{\rm T}^2}{M^2}\,. \qquad (39)
$$

The helicity-two component vanishes as $k_{\rm T}\to 0$, as does the helicity-zero also. However, the helicity-one component ($\propto p_{\rm T}^2$) stays non-zero, in general. In the quark model coupling to two photons, both a_{+-} and a_{+L} are zero [14, 12], and so in this model the cross section will have the same features as that of a 0^{-+} meson, namely a ϕ distribution $\propto \sin^2\phi$, a cross section that vanishes for $k_{\rm T}\to 0$, and helicity-zero dominance.

For central production in hadronic reactions mediated by a vector Pomeron we have to distinguish two cases, namely $a_{+L}\neq 0$ or $=0$. In both cases the helicity-two component is suppressed. In the first case we have a cross section that survives at small $k_{\rm T}$. Moreover, at small $k_{\rm T}$ we expect helicity-one dominance and a flat ϕ distribution. In the second case, i.e. for a suppressed helicity-one amplitude, we predict (i) a vanishing 2^- cross section for $k_{\rm T}\to 0$ (recall, both $q_{1\perp}-q_{2\perp}$ and $\sin^2\phi$ vanish for $k_{\rm T}\to 0$), and, provided $A_{++}\neq 0$, helicity-zero dominance as well as a $\sin^2\phi$ distribution (since $(q_{1\perp}^2-q_{2\perp}^2)^2/M^4$ is smaller at low $k_{\rm T}$ than $\sin^2\phi$).

(iv) J^P=3$^+$ and 1$^-$, 3$^-$

With the help of table 1 and (31) it is straightforward to find the $k_{\rm T}$ and ϕ distributions for the 3^+ states and possible (non-$\rm q\bar{q}$) 1^{-+} and 3^{-+} states.

$$
\begin{aligned}
{\rm d}\sigma[3^+] &\sim \frac{1}{2}\,A_{+-}^2\,(1-\delta_1)\,(1-\delta_2)\,+a_{+L}^2\,\frac{k_{\rm T}^2}{M^2}\\
&\quad + D^2\,a_{++}^2\,(1-\delta_1)\,(1-\delta_2)\,\sin^2\phi\\
{\rm d}\sigma[3^-] &\sim \frac{1}{2}\,D^2\,a_{+-}^2\,(1-\delta_1)\,(1-\delta_2)\,+a_{+L}^2\,\frac{p_{\rm T}^2}{M^2}\\
&\quad + D^2\,a_{++}^2\left(\delta\,\frac{a_{LL}}{a_{++}}-\xi_1\sqrt{(1-\delta_1)\,(1-\delta_2)}\,\cos\phi\right)^2\\
{\rm d}\sigma[1^-] &\sim a_{+L}^2\,\frac{p_{\rm T}^2}{M^2}+D^2\,a_{++}^2\left(\delta\,\frac{a_{LL}}{a_{++}}-\xi_1\sqrt{(1-\delta_1)\,(1-\delta_2)}\,\cos\phi\right)^2\,. \qquad (40)
\end{aligned}
$$

Here we have used that $A_{LL}=D\delta a_{LL}$ for 1^- and 3^- mesons.

As a specific example[6] we illustrate the 1^{--} which can be most immediately relevant in ep $\to$ ep V. Note that Bose symmetry is now not valid and so both the form of the first term in (40) is changed and the factor D^2 is absent in the second term. We find

$$\begin{aligned} d\sigma \;\sim\; & \left(\sqrt{1-\delta_2}\,\frac{a_{L+}}{M}\,\vec{q}_{1\perp} - \eta\,\xi_2\,\sqrt{1-\delta_1}\,\frac{a_{+L}}{M}\,\vec{q}_{2\perp}\right)^2 \\ & + A_{++}^2\left(r - \xi_1\sqrt{(1-\delta_1)\,(1-\delta_2)}\,\cos\phi\right)^2 \;. \end{aligned} \tag{41}$$

In the particular case of forward electroproduction, where $t_2 \to 0$ but $q_{1\perp}^2 = Q^2$ is small, we have approximately

$$d\sigma \sim \frac{Q^2}{M^2}\,a_{L+}^2 + A_{++}^2\,(r-\cos\phi)^2 \;. \tag{42}$$

If for some reason we still have $A_{++} = D\,a_{++}$ or if $A_{++}(Q_1, Q_2 = 0) \sim Q_1^2$ then

$$d\sigma \sim a_{L+}^2 + \frac{Q^2}{M^2}\,(r-\cos\phi)^2 \;. \tag{43}$$

Thus we would expect dominance of transversely-polarized vector-mesons and a longitudinally-polarized component of a characteristic ϕ dependence.

In this section we have given explicit formulae for the CVC case only. While this applies to $e^+e^- \to e^+e^- M$, we have noted that some data involving the Pomeron in proton–proton collisions go beyond this hypothesis. We will discuss elsewhere the detailed phenomenology for both pp and ep-induced reactions.

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