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Evidence that the Pomeron transforms as a non-conserved vector current

Frank E. Close[a]

and

Gerhard A. Schuler[b]

Theoretical Physics Division, CERN, CH-1211 Geneva 23

Abstract

The detailed dependences of central meson production on the azimuthal angle ϕ, t and the meson J^P are shown to be consistent with the hypothesis that the soft Pomeron transforms as a non-conserved vector current. Further tests are proposed. This opens the way for a quantitative description of $q\bar{q}$ and glueball production in $pp \rightarrow p\,M\,p$.

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[a] On leave from Rutherford Appleton Laboratory, Chilton, Didcot, Oxfordshire, OX11 0QX, UK; supported in part by the EEC-TMR Programme, Contract NCT98-0169.

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1 Introduction

In order to understand the dynamics of the proposed glueball filter[1] and to separate glueballs from $q\bar{q}$ states in central production [2] it is necessary to establish the transformation properties of the Pomeron at low momentum transfers. The observations [3] of non-trivial dependence on the azimuthal angle ϕ of the outgoing protons in $pp \to pp + M$ shows that the effective spin of the Pomeron cannot be simply zero [4, 5]. Recently we have shown [4] that many features of the central production of several established $q\bar{q}$ mesons in $pp \to pp + M$, in both the azimuthal, ϕ, and glueball filter, k_T, dependence can be understood if the Pomeron behaves as a conserved vector current. However, there are both empirical [4] and theoretical reasons [5, 6] to believe that this cannot be the whole story and that current **non**-conservation is important, especially at the central meson production-vertex.

In this letter we show that the data clearly require that the Pomeron transforms effectively as a non-conserved vector with a behaviour of a specific type. We propose further tests of this hypothesis and discuss the practicalities of differentiating glueballs from $q\bar{q}$ in $0^{++}, 2^{++}$ meson production.

As in our study of the conserved vector-current (CVC) case [4] we consider the central production of a J^{P+} meson M in high-energy proton–proton scattering. We have shown, with current conservation at the proton–Pomeron vertex (see also [5]), that the cross section may be written as

$$\mathrm{d}\sigma/\mathrm{d}t_1\,\mathrm{d}t_2\,\mathrm{d}\phi\,\mathrm{d}x_\mathrm{F} \sim t_1\,t_2\,(\sigma_2 + \sigma_1 + \sigma_0)\ , \tag{1}$$

where the subscripts i denote the helicity states of the meson and the suppressed pre-factor has the following properties: it is concentrated at $x_\mathrm{F} \approx 0$, steeply falling with decreasing four-momentum transfers, t_i, and finite for t_i approaching $t_i^{\min} \approx 0$ (i.e. approximately proportional to $\exp\{b(t_1 + t_2)\}$).

In the kinematic regime of interest, $|t_i| \ll M^2$ and $x_i \approx (\sqrt{x_\mathrm{F}^2 + 4M^2/s} \pm x_\mathrm{F})/2 \ll 1$, we have

$$\begin{aligned}
\sigma_2 &= \frac{1}{2}A_{+-}^2 \\
\sigma_1 &= A_{+\mathrm{L}}^2 + A_{\mathrm{L}+}^2 - 2\,\eta\,\xi_2\,A_{+\mathrm{L}}\,A_{\mathrm{L}+}\cos\phi \\
\sigma_0 &= (A_{\mathrm{LL}} - \xi_1\,A_{++}\cos\phi)^2 \qquad (\eta = +1) \\
&= A_{++}^2\sin^2\phi \qquad\qquad (\eta = -1)\ ,
\end{aligned} \tag{2}$$

where the subscripts $\pm$ and L refer to the Pomeron helicities, ξ_i are sign factors, and η is the product of the naturality of the meson and the two currents, $\eta = \eta_1\,\eta_2\,\eta_M = \pm 1$. The general structure of the ϕ dependence as a function of J^P is then as follows, from which we will abstract specific tests:

$$\mathrm{d}\sigma[0^-] \sim t_1\,t_2\,A_{++}^2\sin^2\phi \tag{3}$$

$$\begin{aligned}
\mathrm{d}\sigma[0^+] &\sim \left(\sqrt{t_1\,t_2}\,A_{\rm LL} - \xi_1\,\sqrt{t_1\,t_2}\,A_{++}\,\cos\phi\right)^2 && (4)\\
\mathrm{d}\sigma[1^+] &\sim t_1\,t_2\,A_{++}^2\,\sin^2\phi + (\alpha\,\vec{q}_{1\rm T} - \beta\,\vec{q}_{2\rm T})^2 \\
&\qquad (\alpha = q_{2\rm T}\,A_{+\rm L}, \beta = \eta\xi_2\,q_{1\rm T}\,A_{\rm L+} \text{ or } \beta = q_{1\rm T}\,A_{+\rm L}, \alpha = \eta\xi_2\,q_{2\rm T}\,A_{\rm L+}) && (5)\\
\mathrm{d}\sigma[2^+] &\sim \left(\sqrt{t_1\,t_2}\,A_{\rm LL} - \xi_1\,\sqrt{t_1\,t_2}\,A_{++}\,\cos\phi\right)^2 \\
&\qquad + (\alpha\,\vec{q}_{1\rm T} - \beta\,\vec{q}_{2\rm T})^2 + t_1\,t_2\,\frac{1}{2}\,A_{+-}^2 && (6)\\
\mathrm{d}\sigma[2^-] &\sim t_1\,t_2\,A_{++}^2\,\sin^2\phi \\
&\qquad + (\alpha\,\vec{q}_{1\rm T} - \beta\,\vec{q}_{2\rm T})^2 + t_1\,t_2\,\frac{1}{2}\,A_{+-}^2\ . && (7)
\end{aligned}$$

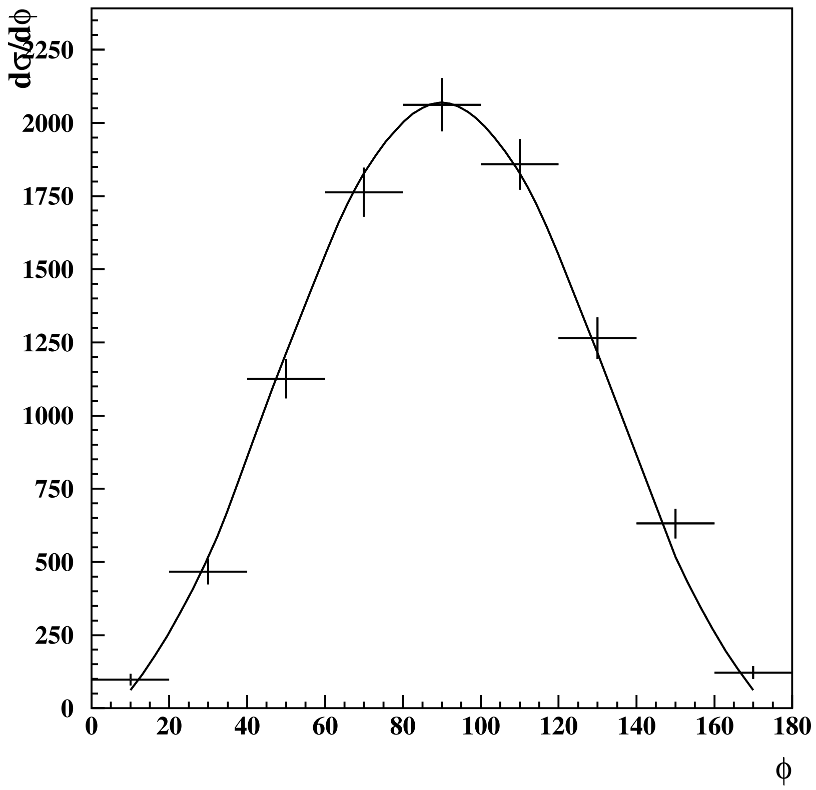

Figure 1: The $\sin^2\phi$ prediction compared to η' production at WA102 [3].

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

The pseudoscalar cross sections behaves as $\mathrm{d}\sigma \sim t_1\,t_2\,A_{++}^2\,\sin^2\phi$. The ϕ dependence is hence independent of the 0^- being a glueball or a quarkonium state [7, 8] and thereby provides an immediate test of the Pomeron transformation properties. The observed η and η' indeed have the $\sin^2\phi$ behaviour (see, for example,

Figure 1). They also exhibit the $t\,e^{bt}$ behaviour [3]. There is the interesting possibility that glueball production might have a compensating $(t_1\,t_2)^{-1}$ Pomeron–Pomeron–glueball vertex that could give a finite cross section as $t_{1,2} \to 0$ and provide a dynamical discrimination.

The $t_1\,t_2$ factor in (3) originates from the fusion of two transversely-polarized currents (TT) and is hence independent of whether the current is conserved or not. The question of current conservation becomes testable if longitudinal polarization (L) can contribute, as in the case of all other states,

The 1^+ state is interesting because Bose symmetry ($A_{++} \propto (t_1 - t_2)^2$ in (5)) suppresses the TT part (with helicity zero) and leaves TL (with helicity one) dominant. As such this becomes sensitive to (non)-conserved vector effects.

In order to exploit current (non)-conservation we consider three scenarios:

i) Current conservation (Model C):

$$q \cdot M = 0\ . \tag{8}$$

ii) Non-conservation (Model B):

$$q \cdot M = O(1)\ , \tag{9}$$

or (Model A):

$$q \cdot M = O(\sqrt{-t})\ . \tag{10}$$

We shall show how data discriminate among these alternatives and provide a consistent solution. In (8–10) q ($t = q^2$) denotes the four-momentum of the current (i.e. $q = q_1$ or q_2). Since the longitudinal polarization vector $\epsilon_{\rm L} \sim q/\sqrt{-t}$ as $t \to 0$ we obtain $\epsilon_{\rm L} \cdot M \sim q \cdot M/\sqrt{-t}$ in contrast to $\epsilon_\pm \cdot M = O(1)$ at small $|t|$. Model C thus corresponds to the conserved-vector hypothesis (CVC): $M_{\rm L}/M_{\rm T} \sim \sqrt{-t}/\mu$, where μ is a mass scale. Model B is what was argued [5] to correspond to the soft ("Donnachie–Landshoff") Pomeron, $M_{\rm L}/M_{\rm T} \sim \mu/\sqrt{-t}$ where $\mu \simeq M$, the mass of the produced meson. Contrary to [5], we anticipate that μ is a rather small mass scale, related to constituent binding or to instanton size. Finally, model A is a possible further alternative where longitudinal and transverse amplitudes have similar strengths, $M_{\rm L}/M_{\rm T} \sim 1$. We now illustrate this in the case of 1^{++}.

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

Here we concentrate on the fusion of two identical currents (i.e. photon-photon or Pomeron-Pomeron). Then σ_0 is as for the pseudoscalars but obeying Bose symmetry

$$t_1\,t_2\,\sigma_0 = t_1\,t_2\,\mu^2\,\hat{A}^2_{++}\,\frac{(t_1 - t_2)^2}{\mu^4}\,\sin^2\phi\ . \tag{11}$$

Here and in the following the hatted quantities denote the residual, dimensionless amplitudes, which, in general, are of order one[1]. As for the pseudoscalars, the

[1] That is we generally assume that $\hat{A}_{\lambda\lambda'} \approx \pm 1$ at small $|t_i|$. Additional dynamics might change this and yield, for example, $\hat{A}_{\rm LL} = c\,k_{\rm T}/\mu$.

helicity-zero part is independent of the model assumption.

Bose symmetry allows us to determine the sign of ξ_2 in σ_1, (2,5), and we find

$$
\begin{aligned}
t_1\, t_2\, \sigma_1 = \quad & \mu^2\, t_1\, t_2\, 4\, \sin^2 \tfrac{\phi}{2}\, \hat{A}_{+L}^2 && \text{model A} \\
& \mu^4\, k_{\rm T}^2\, \hat{A}_{+L}^2 && \text{model B} \\
& t_1\, t_2\, k_{\rm T}^2\, \hat{A}_{+L}^2 && \text{model C}\ . \qquad (12)
\end{aligned}
$$

Recall that $k_{\rm T} \to 0$ implies $\phi \to 0$ but $\phi \to 0$ yields $k_{\rm T} \to 0$ only if $t_2/t_1 \to 1$. We can make this manifest by introducing the following quantities:

$$
\epsilon = \frac{\sqrt{t_1\, t_2}}{\mu^2}\ , \qquad r = \sqrt{\frac{t_2}{t_1}} \qquad (13)
$$

so that

$$
\frac{t_1 - t_2}{\mu^2} = \epsilon\, \frac{1 - r^2}{r}\ , \qquad \frac{k_{\rm T}^2}{\mu^2} = 4\, \epsilon \left(\sin^2 \frac{\phi}{2} + \frac{(1-r)^2}{4\, r} \right)\ . \qquad (14)
$$

This allows us to write the 1^+ cross section as

$$
\sigma \sim \epsilon\, \left(2\, \mu^3\, \hat{A}_{+\rm L} \right)^2\, p_1\, \left\{ \sin^2 \frac{\phi}{2} + p_2 + p_3\, \sin^2 \phi \right\}\ , \qquad (15)
$$

where

Model	Variable p_1	p_2	p_3
A	ϵ	0	$\epsilon^2\, F$
B	1	κ	$\epsilon^3\, F$
C	ϵ^2	κ	$\epsilon\, F$

$$
\kappa = \frac{(1-r)^2}{4\, r}
$$

$$
F = \frac{1}{4} \left(\frac{1 - r^2}{r} \right)^2 \left(\frac{\hat{A}_{++}}{\hat{A}_{+\rm L}} \right)^2\ . \qquad (16)
$$

Note that p_3 characterizes the strength of the helicity-zero part relative to the helicity-one component.

The kinematics of the WA102 experiment is such that in average $r \geq 0.6$, i.e. close to unity. We make the following observations.

- Helicity-one dominance in all cases for integrated cross sections. This is seen in the WA102 data [3]. It is amusing that a dominant helicity-one component was observed already at the ISR 15 years ago [9]. This was regarded as an "unusual feature" and, to the best of our knowledge, has remained a puzzle until now. We predict an enhanced helicity-zero part for ϕ around $\pi/2$ and asymmetric t values (implying small F).
- Vanishing cross sections for $k_{\rm T} \to 0$ in all cases, also in agreement with data [3, 10] (recall that $k_{\rm T} \to 0$ implies $r \to 1$).

- A strong $t_1 t_2$ suppression for model A and an even stronger one for model C, which is not observed in the data. There are, however, indications [11] that data do not simply follow an e^{bt} distribution but exhibit a weak turn-over at small t, precisely as is predicted in model B. Hence we favour model B.
- Owing to the smallness of the helicity-zero part (p_3) we expect a dominant $\sin^2(\phi/2)$ distribution for model A. Indeed, since r is close to unity, this is also the dominant behaviour for models B and C, modulated by an additional isotropic term.

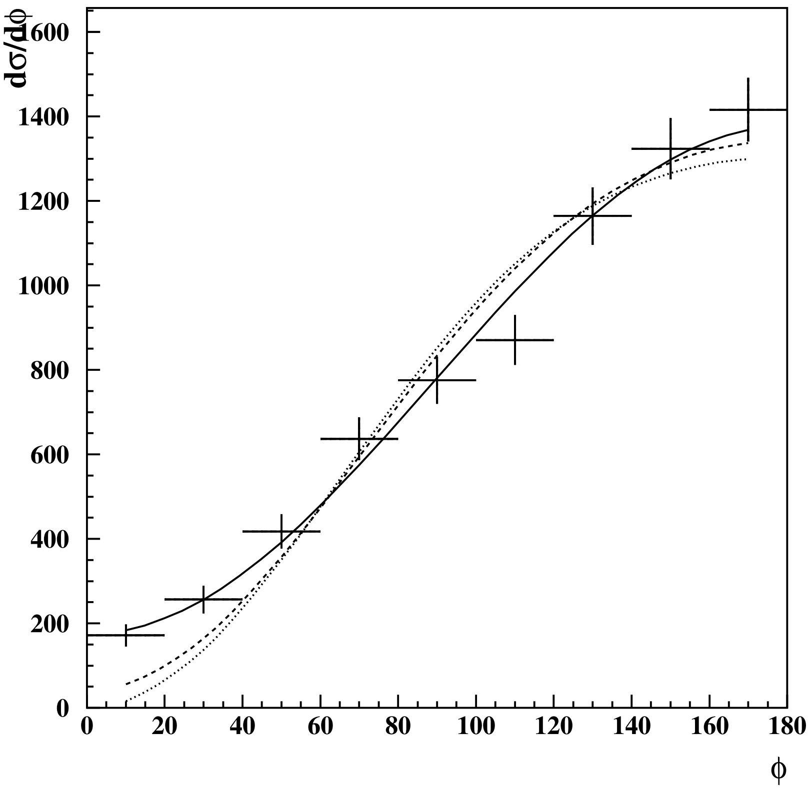

Figure 2: The ϕ distribution of $f_1(1420)$ production at WA102[3] fitted to (15). Solid line: the form of models B and C with p_2 and p_3 as free parameters (fit values are $p_2 = 0.15$ and $p_3 = 5\times10^{-3}$); Dashed line: ditto but with $p_2 = \kappa = 0.03$ fixed to the theoretical value (16) for $r = 0.78$ (fit value is $p_3 = 0.11$); Dotted line: the form of model A, i.e. $p_2 = 0$, with p_3 as free parameter (fit value is $p_3 = 0.15$).

Figure 2 shows a comparsion with the $f_1(1420)$ data from WA102 exhibiting the dominance of the $\sin^2(\phi/2)$ term. An excellent description is described in both model B, or in C, with a small isotropic term ($p_2 = 0.15$) and a very small helicity-zero contribution ($p_3 = 0.005$). If we fix p_2 to its theoretical value using $\langle -t_1 \rangle = 0.145\,\mathrm{GeV}^2$ and $\langle -t_2 \rangle = 0.240\,\mathrm{GeV}^2$, we can determine the mass scale μ by taking $\hat{A}_{++}/\hat{A}_{+\mathrm{L}} = 1$. We find $\mu \approx 0.34\,\mathrm{GeV}$ for $\mathrm{f}_1(1285)$ and $\mu \approx 0.40\,\mathrm{GeV}$ for $\mathrm{f}_1(1420)$ based on the fit values $p_3 = 0.27$ and 0.11, respectively.

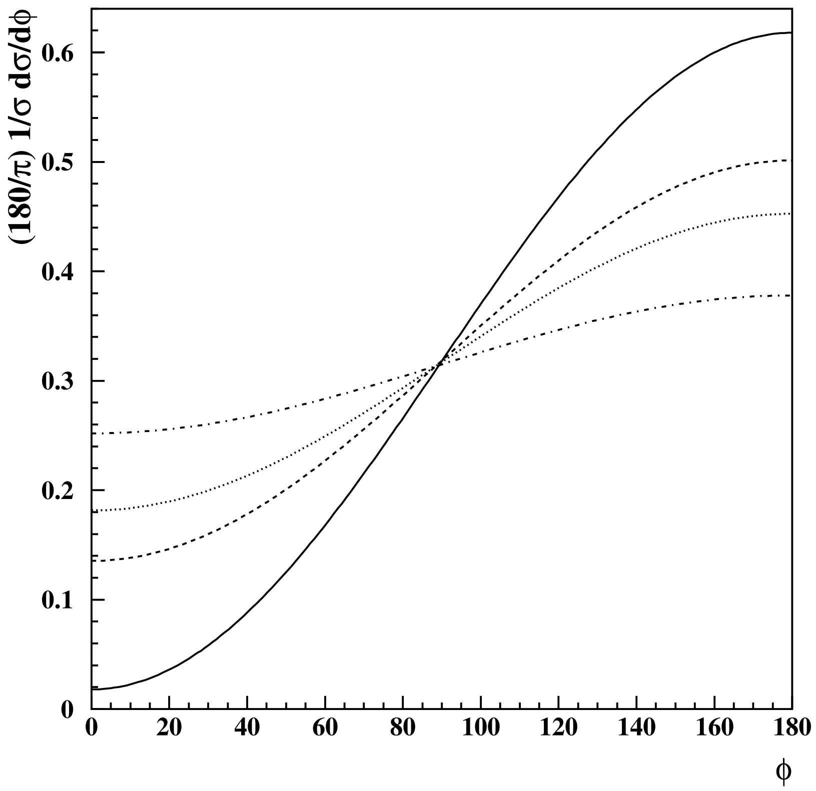

Figure 3: Predictions for the normalized ϕ distribution of 1^+ production for various values of t_2/t_1: 0.5 (solid), 0.1 (dashed), 0.05 (dotted), 0.01 (dashed-dotted). The helicity-zero contribution has been neglected.

To test our prediction further, we propose that experimental data on the ϕ distributions should be analyzed for various r bins. Lowering r from 1 to 0 we predict a change from $\sin^2(\phi/2)$ to a distribution essentially flat in ϕ, see Figure 3.

We now turn to scalars and tensors; these are interesting as both potential glueballs as well as established quarkonia are known to arise. Further it is known that, both in the scalar and tensor sector, there are some mesons that are suppressed at low k_{T} and others that do not show this feature [3]. Therefore it is interesting to separate the Pomeron from the meson dynamics.

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

Obviously $\sigma_2 = \sigma_1 = 0$ but $\sigma_0 = (A_{LL} - \xi_1\, A_{++}\, \cos\phi)^2$ depends on the model:

$$\begin{array}{rll} \mathrm{d}\sigma \sim & \mu^2\, t_1\, t_2\, \hat{A}^2_{++}\, (R - \cos\phi)^2 & \text{model A} \\ & \mu^6\, \hat{A}^2_{\rm LL}\, \left(1 - \frac{\epsilon}{R}\, \cos\phi\right)^2 & \text{model B} \\ & \mu^2\, t_1\, t_2\, A^2_{++}\, (R\,\epsilon - \cos\phi)^2 & \text{model C}\ . \end{array} \tag{17}$$

For all cases,

$$R = \xi_1\, \frac{\hat{A}_{\rm LL}}{\hat{A}_{++}} \tag{18}$$

is, in general[2], a number with absolute value of order one, $|R| = O(1)$. Note that the sign of R cannot be fixed from first principles.

Experimentally the t_i distributions seem to continue to grow at small $|t_i|$ [12] indicating that non-conserving parts must be present in the cross section. In model B we have the possibility to compensate the factor $t_1\, t_2$ in (1) through the $1\sqrt{-t_i}$ enhanced longitudinal amplitude.

Focussing now on model B we note that the ϕ dependence is very sensitive to the ratio $\delta \equiv \epsilon/R = \sqrt{t_1\, t_2}/(R\, \mu^2)$. The ϕ distribution changes from isotropic at small δ to $\cos^2\phi$ at large δ. Data thus allow the determination of the size of μ. The interesting regime is when $|\delta| \approx 1$. If, as suggested by the 1^+ analysis, μ is of the order of Λ_{QCD}, the constituent-quark mass, the average $k_{\rm T}$, or the inverse instanton size, then this would occur for the typical $t_i \sim 0.2$ of WA102. In such a case, depending on whether $\delta = \pm 1$, one expects $\cos^4(\phi/2)$ or $\sin^4(\phi/2)$, see Figure 4.

Data on $f_0(980)$ and $f_0(1500)$ show no suppression at small $k_{\rm T}$ [3]. This implies δ is negative and we predict a dominant $\cos^4(\phi/2)$ dependence. Conversely, the $f_0(2000)$, which is observed to vanish as $k_{\rm T} \to 0$ [3], will have δ positive and hence be maximum as $\phi \to \pi$. To further probe the dynamics of scalar production we propose that the ϕ distributions of the data be analyzed for various bins of $\sqrt{t_1\, t_2}$, see Figure 4.

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

The cross section in the various models is given by

$$\mathrm{d}\sigma^{\rm B} \sim \mu^6\, \hat{A}^2_{\rm LL} \left\{ \left(1 - \frac{\epsilon}{R}\, \cos\phi\right)^2 + \frac{k^2_{\rm T}}{\mu^2} \left(\frac{\hat{A}_{+\rm L}}{\hat{A}_{\rm LL}}\right)^2 + \frac{1}{2}\, \epsilon^2 \left(\frac{\hat{A}_{+-}}{\hat{A}_{\rm LL}}\right)^2 \right\} \tag{19}$$

$$\mathrm{d}\sigma^{\rm A} \sim \mu^2\, t_1\, t_2 \left\{ \hat{A}^2_{++}\, (R - \cos\phi)^2 + \hat{A}^2_{+\rm L}\, 4\, \sin^2\frac{\phi}{2} + \frac{1}{2}\, \hat{A}^2_{+-} \right\} \tag{20}$$

$$\mathrm{d}\sigma^{\rm C} \sim \mu^2\, t_1\, t_2 \left\{ \hat{A}^2_{++}\, (\epsilon\, R - \cos\phi)^2 + \frac{k^2_{\rm T}}{\mu^2}\, \hat{A}^2_{+\rm L} + \frac{1}{2}\, \hat{A}^2_{+-} \right\}\ . \tag{21}$$

[2]See footnote 1.

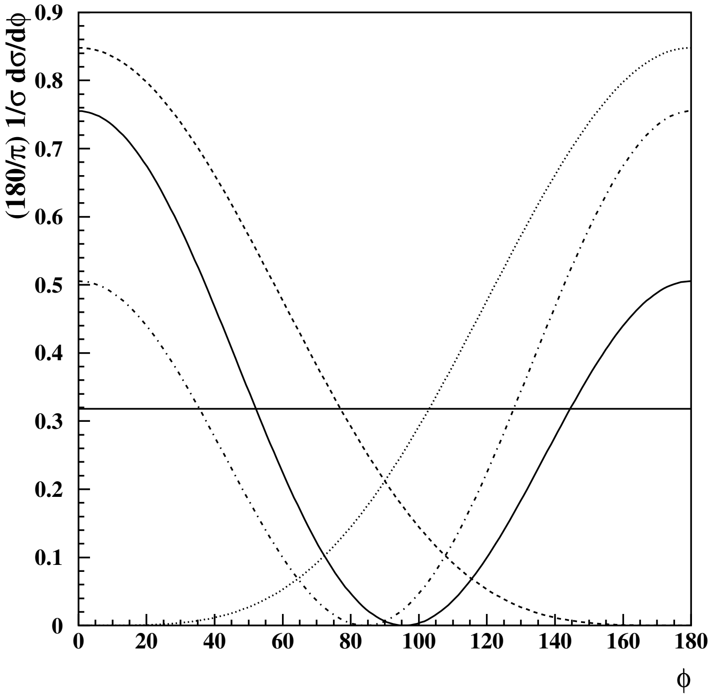

Figure 4: Predictions for the normalized ϕ distribution of 0^+ production in model B, $\mathrm{d}\sigma/\mathrm{d}\phi \sim (1 - \delta\,\cos\phi)^2$, for various values of $\delta = \epsilon/R$: -10 (solid), -1 (dashed), 0 (solid), $+1$ (dotted), $+10$ (dashed-dotted).

The three terms for each case correspond to meson helicity zero, one, and two, respectively. Here we have again made use of Bose symmetry to fix the sign in σ_1.

In CVC, the residual amplitude $\hat{A}_{+-}$ is naturally of order one; hence a large helicity-two component is expected for quarkonium states. The central-production data [3] do not agree with this even though helicity-two dominance is well established and understood for e^+e^- [13, 14]. This marked difference was commented on in our earlier paper and is another motivation for non-CVC dynamics.

As for scalar-meson production, experimental hints [12] for a continuous growth of the t distributions at small $|t|$ favour the current non-conserving alternative B. Let us see whether we can find a consistent picture.

Well-established quarkonia are known to be suppressed at small k_{T} [3]. Referring to (19) this implies that

$$\epsilon^2\,(\hat{A}_{+-}/\hat{A}_{\mathrm{LL}})^2 \ll 1 \qquad \text{and} \qquad \delta = \epsilon/R \approx +1\ . \tag{22}$$

If this was the case we would predict a helicity hierarchy, namely $\sigma_0 \gg \sigma_1 \gg \sigma_2$. This looks consistent since preliminary WA102 data on $f_2(1270)$ and $f_2'(1525)$ [15] seem to support this ordering. Moreover, the condition $|\delta| \approx 1$ is also what we already found from our analysis of scalars.

If $\delta \approx +1$ (in accord with the suppression as $k_{\rm T} \to 0$) then it is sensible to expand in powers of $\sin^2(\phi/2)$

$$\mathrm{d}\sigma^{\rm B} \sim 4\,\mu^6\,\hat{A}^2_{\rm LL}\,\left\{p_2 \sin^4\frac{\phi}{2} + p_1 \sin^2\frac{\phi}{2} + p_0\right\}\ , \tag{23}$$

where

$$\begin{aligned} p_2 &= \delta^2 \\ p_1 &= (1-\delta) + \epsilon\left(\frac{\hat{A}_{+\rm L}}{\hat{A}_{\rm LL}}\right)^2 \\ p_0 &= (1-\delta)^2 + \frac{\epsilon}{4}\,\frac{(1-r)^2}{r}\left(\frac{\hat{A}_{+\rm L}}{\hat{A}_{\rm LL}}\right)^2 + \frac{\epsilon^2}{8}\left(\frac{\hat{A}_{+}}{\hat{A}_{\rm LL}}\right)^2\ . \end{aligned} \tag{24}$$

We see that helicity zero contributes to all p_i, helicity one to p_1 and p_0, and helicity two to p_0 only. Hence we predict that both the helicity structure and the ϕ dependence should vary with t.

If, as has been suggested elsewhere [1], 2^+ glueballs survive as $k_{\rm T} \to 0$, then we would expect a ϕ distribution more similar to $\mathrm{f}_0(980, 1500)$ than to $\mathrm{f}_2(1270)$.

(v) J^P=2^-

The cross section in model B is given by

$$\mathrm{d}\sigma[2^-] \sim t_1\,t_2\,\mu^2 \hat{A}^2_{++}\,\sin^2\phi + \mu^4\,p_{\rm T}^2 \hat{A}^2_{+\rm L} + \frac{t_1\,t_2\,(t_1-t_2)^2}{2\,\mu^2}\hat{A}^2_{+-}\ , \tag{25}$$

where the three terms on the rhs correspond to helicity zero, one, and two, respectively. The expression for models A and C are given by replacing the helicity-one term by (12) with the substitutions $k_{\rm T} \mapsto p_{\rm T}$ and $\sin(\phi/2) \mapsto \cos(\phi/2)$. We can write (25) as

$$\begin{aligned} \mathrm{d}\sigma[2^-] \ \sim\ \epsilon\,\mu^6 & \left\{\epsilon\,\hat{A}^2_{++}\,\sin^2\phi + 4\,\hat{A}^2_{+\rm L}\left[\cos^2\frac{\phi}{2} + \frac{(1-r)^2}{4\,r}\right]\right. \\ & \left. + \frac{\epsilon^2}{2}\left(\frac{1-r^2}{r}\right)^2\,\hat{A}^2_{+-}\right\}\ . \end{aligned} \tag{26}$$

Experimentally the 2^- states are known to be suppressed at small $k_{\rm T}$ [3]. This implies that the reduced helicity-one amplitude $\hat{A}_{+\rm L}(k_{\rm T} \to 0)$ is suppressed. This happens in the non-relativistic quark model coupling to two photons, where $\hat{A}_{+\rm L}$ is identically zero. If this is also true in the QCD case, then we predict:

- Cross sections that behave as $t\,e^{b\,t}$ at small t, similar to the 0^- ones.
- A small helicity-one contribution.

- A helicity-zero contribution that behaves as $\sin^2\phi$.
- A helicity-two contribution that is isotropic in ϕ.
- A ratio σ_2/σ_0 that vanishes for $t_2 = t_1$ and increases with increasing difference $|t_2 - t_1|$.

In summary, we eagerly await new results on the ϕ and t dependences for central meson production enabling the parameters, as described in this paper, to be determined. Once these parameters are determined, the dynamical nature of the mesons will become clear.

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