Abstract: We discuss here a quark model, which explicitly embodies quark confinement, and show how one can trace all processes leading to multi-hadron final states, to one and the same basic underlying mechanism, involving the annihilation of an energetic $q\bar{q}$ system. This picture leads to jets and multi-hadron final states always with a very similar structure, except for the kinematic difference between the various inducing processes.

Résumé: Nous discutons ici un modèle de quark qui garantit explicitement l'emprisonnement des quarks, et nous montrons comment tous les processus menant à la production d'états multi-hadroniques peuvent être reliés au même mécanisme fondamental qui comprend l'annihilation du système $q\bar{q}$ très énergétique. Cette image conduit à des "jets" et à des états finaux multi-hadroniques ayant toujours une structure très semblable, sauf pour les différences cinématiques entre les divers processus de départ.
1. MOTIVATION AND OBJECTIVES

At one time or another almost all aspects of hadronic physics have been discussed in terms of quarks. This includes

- the spectrum of hadrons and their static properties;
- weak and electromagnetic interactions;
- elastic scattering of hadrons;
- large $p_T$ phenomenon.

However, quarks or states with quark quantum numbers have never been seen in hadronic final states, so in quark models one tends to leave the question of hadronic final states in a rather unsatisfactory state. The present talk is concerned with just this aspect of hadronic physics within a quark framework. We shall use a simple parametrization of quark confinement and try to make some sense of hadronic final states emerging from such a framework. Our basic premise will be that all multi-hadron production processes can be traced back to one and the same underlying mechanism.

We may remark that even if in some extreme asymptotic domain quarks are produced, to the extent that they are not seen in sub-asymptotic regions, our considerations should apply as long as we still insist on describing processes in terms of quark lines. The question of the observability of quarks is of course a fundamental one and the answer will, in an essential way, dictate how we should describe quark degrees of freedom. For example, if quarks can be produced, then we can with some confidence assume they can be described in terms of a conventional quantum field theory. If, however, the converse is true, it is by no means clear that this can be done, since the meaning of quark degrees of freedom no longer follows the conventional wisdom and the quark notion may at most be a convenient way of parametrizing the internal symmetry and structure of hadrons. We therefore feel that it is important to unequivocally establish the utility of the quark notion in hadronic physics. One might reasonably argue\(^{1-4}\) that this has already been adequately demonstrated, provided one ignores the question of the final state. However, it is our view that we can no longer ignore the latter problem and we shall devote our discussion solely to it. We do this by addressing our considerations to the following empirical facts which, although admittedly may be interpreted in an over-simplified way, allow us to concentrate on what we see as the essential clues to the quark-hadronic puzzle.
a) Early scaling in all processes involving multi-hadron final states
\[ \sigma_{\text{tot}}^{\gamma \gamma \rightarrow X(s)} \sim s^2 E^\varepsilon \] with \( \varepsilon > 0 \)
\[ \sigma_{\text{tot}}^{\gamma \gamma \rightarrow X(q^2, s)} \sim F(q^2/s) \]
\[ \sigma_{\text{tot}}^{e^+e^- \rightarrow X(q^2)} \sim \sigma_{e^+e^- \rightarrow \mu^+\mu^- R}. \]

b) The apparent universality in the structure of hadronic final states, after allowing for kinematic differences
\[ n_{\text{ch}} \sim m (c \log s) \quad m = 1, 2; \quad c = 1 \]
\[ P_T \sim 300 \text{ MeV} \]
short range correlations
fast particle production, even in rare events at large angles, always occurs in jets with a very similar structure.

c) The apparent geometrical properties of hadronic amplitudes.

d) The Pomeron appears to be a pole with intercept above unity even up to the highest available energies.

Let us end this introductory discussion by recalling the naive quark-parton ideas\(^1, \, ^2\) and the massive quark model\(^3\) and their application to hadron-hadron scattering and deep inelastic current induced processes. A long time ago it was thought that the approximate 2/3 ratio of the meson-baryon and baryon-baryon total cross-sections could be understood in terms of the additive quark model (Fig. 1a), while more recently the scaling laws in \( \gamma \gamma \rightarrow X \) and \( e^+e^- \rightarrow \gamma \gamma \rightarrow X \) could be understood in terms of the quark-parton model, Fig. 1b and Fig. 1c, respectively

![Diagram](image)

**Fig. 1** The quark-parton diagrams

The dashed lines in Fig. 1 indicate we are interested in the discontinuities of these diagrams; Fig. 1a leads to \( \sigma_{\text{tot}}^{\gamma \gamma \rightarrow X} = n a n b q q' \), while Figs. 1b and 1c lead to the scaling laws indicated in (a) above. However, it is
difficult to reconcile these mechanisms with the absence of quark final states and one possibility is to replace them by the diagrams of Fig. 2a, b and c.

![Diagrams](image)

Fig. 2 The corresponding massive quark model (MQM) diagrams

The desired scaling properties result from the diagrams in Fig. 2 if we assume

1) The \( q\bar{q} \) Green's function has a simple scaling behaviour

\[
G_{qq}(s,t,q_1^2,q_2^2,q_3^2,q_4^2) = s f(t,q_1^2,q_2^2,q_3^2,q_4^2).
\]

2) There are no singularities in quark lines. Furthermore, this must be achieved with a universal peaking in the quark four-momentum squared variables \( q_i^2 \) around some small value \( q_i^2 = m^2 \).

The ansatz (1) and (2) are the main ingredients of the MQM of Preparata\(^5\). In such a model the basic mechanism for particle production is the annihilation of excited \( q\bar{q} \) states (Fig. 3).

![Diagrams](image)

We notice that the MQM relates the scaling behaviour \( o_{\text{tot}}^{X}(s) = \text{constant} \) to the deep inelastic scaling laws and indeed after a \( t \)-channel unitarization of Fig. 2a, leads to a \( s^{1+\epsilon} \) Pomeron\(^6\). However, the question of the structure of the final states remains unanswered and we must find a suitable space-time description of confined quarks.

2. BASIC INGREDIENTS OF A SPACE-TIME DESCRIPTION OF CONFINED QUARKS

We sketch here the basic ideas behind a simple parametrization of quark confinement proposed by Preparata and myself\(^7,8\) which, although in spirit is similar to the MIT bag\(^9\), it is designed so as to allow hadronic scattering and production processes to be discussed in a detailed way. The essential
difference with the MIT bag lies in the simple way we realize translational invariance which, like the MQM, allows one to write down a graphical computational scheme closely analogous to the quark-parton model\(^2\).

2.1 Quark confinement

In any model involving quark lines one deals with Green's functions involving \( n \) quark (antiquark) legs and \( N \) hadron legs, where \( n = 2, 3, 4, 6, \ldots \) and \( N = 0, 1, 2, \ldots \).

\begin{figure}
\centering
\includegraphics[width=0.5\textwidth]{quark_confinement.png}
\caption{Fig. 4}
\end{figure}

Let us consider such an object as a function of the quark space-time coordinates \( \{x_i\} \) and the hadronic momentum \( \{p_i\} \). Then the simplest statement of quark confinement (for scalar quarks) is to assume

\[(\Box_i + m_i^2) \tilde{G}(x_1, x_2, \ldots, x_n; p_1, \ldots, p_N) = 0 \quad i = 1, n\]

for \( y_{ij} = x_i - x_j \in R_{ij}(p_1, \ldots, p_N) \) \( j = 1, n \)

and \( \tilde{G} = 0 \) for \( y_{ij} \notin R_{ij} \),

where \( R_{ij} \) is a compact space-time domain or cell. These equations describe quarks freely moving inside some compact space-time region dependent on the configuration of hadrons \( \{p_i\} \). One can easily convince oneself that these quark-hadron Green's functions have the following three important properties:

i) in momentum space \( G(q_1, q_2, \ldots; p_1, \ldots, p_N) \) is an entire function in the variables \( \{q_i^2\} \) which are strongly peaked around \( q_i^2 = m^2 \) (depending on the size of the confinement region);

ii) implies a spectrum of hadron states;

iii) the statement is compatible with translational invariance, the implementation of which leads to the very desirable conclusion that these Green's functions conserve momentum, i.e. \( \sum_i q_i = \sum_i p_i \).

The simplest dynamical scheme one can envisage is one in which the higher order Green's functions are built up of the lower order ones in much the same way as the perturbation expansion of quantum electrodynamics (QED). The simplest objects will then be
2.2 By studying the $qq$ meson wave function, which in configuration space is depicted by

$$\psi(x_1, x_2, p) = \psi(p)$$

while in momentum space, after separating out the four-momentum conservation $\delta$-function, we depict it by

$$q_1 = \frac{1}{2} p \cdot q$$

and requiring the peaking in $q_1^2$ and $q_2^2$ to be independent of $p^2$, we obtain for a spherical boundary condition in the rest frame of $p^\mu$ the spectrum

$$M^2_{n,J} \approx \frac{\pi}{R^2} (2n + J)$$

2.3 Furthermore, by defining a charge normalization condition of the form

$$\int |\chi(q,p)|^2 d^4q = Z$$

and allowing only states with $J \leq MR_0/2$ to couple to the $qq$ system, we obtain a $qq$ four-point function with the following discontinuity

$$\text{disc}_s G(s,t,q_1^2,q_2^2,q_3^2,q_4^2) = s \sigma_{qq} F_{RT}(t) \Delta_{R^2}(q_1^2,q_2^2) \Delta_{R^2}(q_3^2,q_4^2)$$

with

$$F_{RT}(t) = 2J_1(R\sqrt{-t})/R_1\sqrt{-t}$$

and

$$\Delta_{R^2}(\sigma_1,\sigma_2) \approx \delta_{R^2}(\sigma_1 - m^2) \delta_{R^2}(\sigma_2 - m^2); \quad \delta_{R^2}(\sigma) = \frac{1}{\pi} \frac{\sin R^2\sigma}{\sigma}.$$ 

2.4 Hadronic interaction

The hadrons are envisaged to interact only through the space-time overlap of their quark wave functions and the lowest order vertex is the triangular graph.
where \((\mu^2)^3\) is an effective coupling constant which turns out to be small.

Since we assume translational invariance of all Green's functions, the above graph can be computed in much the same way as a Feynman graph. However, since the vertices are entire functions \(q_i^2\) (the quark mass square variables), the integrals run over a region in which these variables are kept real. A possible complete coupling scheme is discussed in Ref. 7, where it is demonstrated that these notions lead to a sensible model of strong interactions. We refer to Ref. 7 for details.

3. PROPERTIES OF SUPER-EXCITED \(q\bar{q}\) STATES

3.1 Asymptotic width

\[ \gamma = \lim_{N \to \infty} M^2 N \approx \frac{2\pi}{R^2} = 2\Delta M^2, \]

where \(\Delta M^2\) is the level spacing. This means the levels totally overlap locally, so that in effect we have the \(q\bar{q}\) system coupling to a hadronic continuum.

3.2 Super-excited \(q\bar{q}\) system (fire cylinderoid)

If we consider the states that build up the \(q\bar{q}\) four-point function as \(s = M^2 \to \infty\), then choosing a polarization axis \(z\) and summing over \(J\) (\(0 \leq J \leq \frac{M_T}{2}\)) all degenerate states of mass \(M\), we see that the quarks are effectively propagating in the cylindrical region shown in Fig. 6.

3.3 Chain decay mechanism of super-excited \(q\bar{q}\) states

By studying the basic three-meson vertex we find that these highly excited \(q\bar{q}\) states decay preferentially by cascading down in energy emitting the highest \(q\bar{q}\) states \((\pi, K, \rho, \ldots)\) as quanta. This leads to the following chain decay mechanism (Fig. 7), in which the intermediate states run through all \(J\) compatible with \(0 \leq J \leq \frac{M_T}{2}\).
Fig. 7 Cascade decay mechanism of super-excited $q\bar{q}$ states

One can see that the above mechanism leads to a longitudinal decay distribution by picturing the optimal overlap of the highly excited $q\bar{q}$ states (i.e. fire cylindroids) (Fig. 8a). Clearly one fire cylindroid prefers to go into another one lying concentrically within it and having the pion running along the $z = t$ axis. This results in a longitudinal decay distribution (Fig. 8b).

![Diagram of cascade decay mechanism](image)

(a) Optimal overlap of $q\bar{q}$ fire cylindroids
(b) Longitudinal decay distribution

Fig. 8

4. DYNAMICAL EQUATIONS GOVERNING THE HADRONIC FINAL STATES EMERGING FROM SUPER-EXCITED $q\bar{q}$ STATES

In order to determine the nature of the hadronic final states emerging from the annihilation of an energetic $q\bar{q}$ pair, we cast the chain decay mechanism in the form of the following set of integral equations (see Ref. 8 for details).

$$
\begin{align*}
q & \rightarrow G \rightarrow \pi \\
\bar{q} & \rightarrow G \rightarrow \pi
\end{align*}
$$

(1)

$$
\begin{align*}
q & \rightarrow G \rightarrow G_6 \\
\bar{q} & \rightarrow G \rightarrow G_6
\end{align*}
$$

(II)

$$
\begin{align*}
q & \rightarrow G_9 \rightarrow 2G_6 \\
\bar{q} & \rightarrow G_9 \rightarrow 2G_6
\end{align*}
$$

(III)

(1 + 2 + 3)
Using the variables $k'_i = (k_{T1}', x_1/\sqrt{s}/2)$, the first equation (I) leads to the eigenvalue condition

$$\left[ 1 - \frac{\lambda}{4\pi} \int_0^1 \frac{dx'}{x'} \left( 1 - x' \right) \int \frac{d^2k_T'}{(2\pi)^2} K[x', k_T'] \right] \sigma \rho_q = 0 \quad (I')$$

while the second equation (II), which governs the one-particle distribution $q\overline{q} \to \pi(k_1) + X$ has the form

$$C_{q\overline{q}}[s, x, k_T] = \frac{\lambda}{2(2\pi)^3} s(1 - x) \sigma \rho_q K[x', k_T'] +$$

$$+ \frac{\lambda}{4\pi} \int_0^1 \frac{dx'}{x'} \int \frac{d^2k_T'}{(2\pi)^2} K[x', k_T'] C_{q\overline{q}}[1 - x, s \frac{x'}{\sqrt{1 - x}} k_T - k_T']$$

If we define the scaling function

$$F(x) = \int \frac{d^2k_T'}{(2\pi)^2} \frac{1}{s} C_{q\overline{q}}[s, x, k_T'] ,$$

then from (II'a) we see that it satisfies the equation

$$F(x) = F_0(x) + \frac{\lambda}{4\pi} \int_0^1 \frac{dx'}{x'} (1 - x') K(x') F \left( \frac{x'}{\sqrt{1 - x}} \right) \quad (II'b)$$

with

$$K(x') = \int \frac{d^2k_T'}{(2\pi)^2} K[x', k_T'] .$$

Equation (II'b) can be solved by taking its Mellin transform defined by

$$\tilde{F}(\omega) = \int_0^1 dx / x^{1+\omega} F(x) ,$$

with the result

$$F(x) = \frac{i^{\omega+c}}{2\pi i} \int_{i\infty+c} \frac{d\omega}{\omega} \frac{1}{\left( \frac{1}{\lambda} - \tilde{K}(\omega) \right)} , \quad (II'c)$$

where

$$\tilde{K}(\omega) = \frac{1}{4\pi} \int_0^1 \frac{dx'}{x'} (1 - x')^{1+\omega/2} \int \frac{d^2k_T'}{(2\pi)^2} K[x', k_T'] . \quad (II'd)$$

Comparing (II'd) with eigenvalue condition (I'), we see that $1 - \lambda \tilde{K}(0) = 0$ and that the leading behaviour as $x \to 0$ is determined by a pole at $\omega = 0$ in the denominator of (II'c). From this it follows $F(x) \to \text{constant}$ as $x \to 0$.

The analysis can be readily extended to the region $x \approx 1$ and to the twice inclusive distribution $q\overline{q} \to \pi(k_1) + \pi(k_2) + X$. The results of such an analysis are summarized below.

i) $F(x) = \frac{1}{C} \int_{q\overline{q}} \frac{d\sigma_{q\overline{q}}}{dy} \left. \frac{y}{1/2Y} \right|_{y<<1/2Y} C$ (i.e. we have a plateau in rapidity).

ii) The mean multiplicity $n_{q\overline{q}} = CY, Y = 1/2 \log (s/s_0)$. 

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iii) $F(x) \xrightarrow{(x+1)^{-\lambda}} K(0)(1-x)^\gamma$

(i.e. no quark fragmentation). ($\gamma = 1$ for scalar quarks and $\gamma = 2$ for spin 1/2 quarks)

iv) $\langle k(T) \rangle (x) \xrightarrow{x+0} f(n); \langle k^2 \rangle = \pi/R^2 \quad$ for $x \sim 0$.

v) $\frac{1}{\sigma_{qq}} \frac{d^2 \sigma}{dy_1 dy_2} |y_1, |y_2| < 1/2 \gamma 2C^2 \langle n(n-1) \rangle = C^2 \gamma^2$

(i.e. $D^2 = \langle n^2 \rangle - \langle n \rangle^2 = \langle n \rangle$)

(i.e. we have a Poisson multiplicity distribution).

vi) $C(y_1 - y_2) = \int dy_1 dy_2 \delta(y_1' - y_2' - y_1 + y_2) \times$

$\times \left\{ \frac{1}{\sigma_{qq}} \frac{d^2 \sigma}{dy_1 dy_2} - \frac{1}{\sigma_{qq}} \frac{d\sigma}{dy_1} \frac{1}{\sigma_{qq}} \frac{d\sigma}{dy_2} \right\}$

$\sim e^{-p|y_1 - y_2|}$ with $p \sim 1 - 2$

(i.e. we have short range correlations in rapidity).

vii) The transverse momentum is not locally compensated in rapidity in contrast to multiperipheral dynamics.

5. COMPARISON WITH MULTIPERIPHERAL DYNAMICS

The lifetime of the super-excited $qq$ states is given approximately by $\tau = 1/\Gamma \sim R^2 \sqrt{s}/2\pi$, while the period of quark trajectory is $T = R^2 \sqrt{s}$. This means the excited $qq$ system decays many times during one orbit, so, in effect, the quarks behave like free particles bremsstrahlung pions and we have the following physical picture, which we compare with a multiperipheral mechanism.

![Fig. 9 Decay of a qq fire cylindroid](image)

![Fig. 10 Corresponding picture of a MPM in c.m.s.](image)
In both cases the final states look very similar with \( n \sim \log s \) and \( \langle k_T \rangle \) fixed. However, the dynamics of the fire cylindroid leads to a fixed pole in the q\( \bar{q} \) Green's function, while the multiperipheral mechanism leads to a moving Regge pole.

6. CONSEQUENCE FOR CURRENT INDUCED AND HADRONIC PRODUCTION PROCESSES

i) \( e^+e^- \rightarrow \text{hadrons} \): Involves the excitation of a simple q\( \bar{q} \) cylindroid and the processes can be pictured in the following way

\[
\tilde{n}_{e^+e^-} = c \log S/So
\]

\[
\langle k_T \rangle_{e^+e^-} = \langle k_T \rangle_{q\bar{q}}
\]

(the transverse momentum being defined with respect to the jet axis in each event).

ii) \( hh' \rightarrow \text{hadrons} \): Involves the excitation of one or more pairs of q\( \bar{q} \) fine cylindroids, each pair totally overlapping in phase space. This can be pictured in the following way.

\[
\tilde{n}_c = n_{qq} - 2c \log S/So
\]

\[
\langle k_\perp \rangle_c = \langle k_\perp \rangle_{q\bar{q}}
\]

One can draw similar pictures for \( \gamma\nu \pi \rightarrow X \), which will, in principle, involve excitations of both types (i) and (ii) corresponding to the current and the hadronic plateaus, respectively, which at very large energies should have different heights.

7. THE POMERANCHUCK MECHANISM

The lowest order term in hadron-hadron scattering corresponding to mechanism (ii) above leads to the graph shown earlier in Fig. 2 and corresponds to a fixed pole at \( j = 1 \). Such a Born term violates by itself t-channel unitarity and it should be iterated in the t-channel (Fig. 11).
In Ref. 6, using techniques proposed by Gribov\textsuperscript{10}, we showed that such an iteration leads to the following Dyson equation for the Pomeron propagator.

\[ D(j,t) = D_0(j,t) + D_0(j,t) \cdot V(j,t) \cdot D(j,t), \]

i.e.

\[ D(j,t) = \left[ D(j,t) - V(j,t) \right]^{-1}, \]

where \( V(j,t) \) is the Reggeon potential generated by the two-pion by intermediate states in the t-channel, which seem to have the largest effect. The parameters obtained in Ref. '6 for the Pomeron trajectory are given by

\[ \alpha_p(t) = 1 + \varepsilon + \alpha_p' \cdot t + \alpha_p'' \cdot t^2 + \ldots, \]

with \( \varepsilon = 0.06, \alpha_p' = 0.25, \) and \( \alpha_p'' = 0.5. \)

Clearly such a Pomeron cannot be the whole story, since it must at least be unitarized by s-channel iterations. Further, a bare Pomeron pole with intercept above a certain critical value may be irreconcilable with the inelastic t-channel unitarity constraints. However, it has been argued by Amati, Caneschi and Jengo\textsuperscript{11} that a Pomeron with intercept slightly above unity can have a perturbative s-channel expansion up to some quite high energy, after which a non-perturbative behaviour and t-channel effects take over. We shall appeal to the latter possibility and briefly discuss the final state structure corresponding to the Pomeron mechanism discussed above.

One can readily convince oneself that the Born term, which corresponds to the overlap of two \( q \bar{q} \) clusters of particles, each of which having a Poisson multiplicity distribution and a plateau in rapidity, also has a Poisson multiplicity distribution, but with twice the mean multiplicity and a plateau of twice the height.

As the energy increases the picture begins to look more like that shown in Fig. 12.
In the weak coupling solution in Ref. 6 the probability distribution of the giant clusters is given by

\[ P_{cl}(N, N) = \frac{(\bar{N} - 1)^{N-1}}{(N - 1)!} e^{-(\bar{N}-1)}, \]

with \( \bar{N} = \varepsilon \log s/s_1 + 1 \). Over ISR energies \( \bar{N} = 1.5-2 \), taking \( s_1 = 6 \text{ GeV}^2 \).

If one calculates the one-particle distribution corresponding to Fig. 12, one obtains the interesting result that the height of the central plateau \( h = (1/s)(d\sigma/dy) \bigg|_{y=0} \) rises with energy like the total cross-section. Furthermore, although we have particles strung out in rapidity, the transverse momentum is not compensated locally in rapidity. These features seem to be borne out by the experimental data\(^{12}\).

8. GEOMETRICAL PROPERTIES OF HADRONIC AMPLITUDES

One of the attractive features of the naive quark model is that it provided a simple interpretation of geometrical properties such as the Chou-Yang conjecture\(^{13}\), as long as one ignores the final state problem. In taking this problem into account we have been lead some way from the simple geometrical quark picture. However, by virtue of the scaling property of qq Green's function, its strong peaking in the quark mass squared variables, and its special t-dependence, one can go some way towards recovering the geometrical properties. These properties lead to the following approximation\(^5\) (illustrated here graphically).

For example, in \( e^+e^- \rightarrow \text{hadrons} \) and \( \gamma_j p \rightarrow X \), respectively, we have

\[ = \]

The constant \( c \) is intimately tied to the normalization of the wave functions and although it can be equal to unity in general, it appears not to be and, in fact, appears somewhat larger. If we now turn to baryon-baryon elastic scattering, we envisage the mechanism in Fig. 13a.
The above approximate identities plus the relation
\[ \int d^2q' F_{R_T^{-q'}} F_{R_T^{-q'}-q'} = F_{R_T^{-q}} \]
leads us to the approximation shown in Fig. 13b. However, the validity of this approximation requires a detailed study of the baryon problem, which we suspect will lead us to the desired geometrical properties.

9. LARGE \( p_T \) RARE OVERLAPS

In addition to the normal events, we can also have rare events arising from large angle overlap of the \( q\bar{q} \) fire cylindroids, which is illustrated in Fig. 14a. When seen in the \( q\bar{q} \) rest system it leads to the final state shown in Fig. 14b. Since in general the \( q\bar{q} \) system will have a component of momentum along the \( z \) axis, so the decay distribution will have the two-jet appearance illustrated in Fig. 14c.

One should allow for this rare overlap to occur after \( n \) steps in the normal mode with \( n = 0, 1, 2, \ldots \), and sum over the various possibilities (Fig. 15a). The incoherent summation can lead to an increase in the associated multiplicity with a large \( p_T \) trigger.
For pp scattering the full dynamical mechanism leading to a large $p_T$ event is shown in Fig. 15b. Monte Carlo calculations based on this mechanism by Preparata and Rossi give encouraging results¹⁴).
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