Resonance-reggeon and parton-hadron duality in strong interactions

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Abstract. By using the concept of duality between direct channel resonances and Regge exchanges we relate the small- and large-$x$ behavior of the structure functions. We show that even a single resonance exhibits Bjorken scaling at large $Q^2$.

PACS. 12.40.Nn Regge theory, duality, absorptive/optical models – 13.60.Hb Total and inclusive cross sections (including deep-inelastic processes) – 11.55.Bq Analytic properties of S matrix – 11.55.Hx Sum rules

1 Introduction

Inspired by recent experimental measurements of the nucleon structure functions at the JLab (CEBAF), we suggest a unified "two-dimensionally dual" picture of the strong interaction connecting low- and high energies (Veneziano, or resonance-reggeon duality with low- and high virtualities ($Q^2$) (Bloom-Gilman, or hadron-parton duality). The basic idea of the unification is the use of $Q^2$-dependent dual amplitudes, employing nonlinear complex Regge trajectories providing an imaginary part of the scattering amplitude related to the total cross section and structure functions, thus saturating the duality by a finite number of resonances lying on the (limited) real part of the Regge trajectories.
The resulting object, a deeply virtual scattering amplitude, $A(s, t, Q^2)$, is a function of three variables, reducing to the nuclear structure function (SF), when $t = 0$, and on-shell hadronic scattering amplitude for $Q^2 = m^2$. It closes the circle in Fig. 1. We use this amplitude to describe the background as well as the resonance component [7].

The $Q^2$ dependence of the residue functions here will be chosen in such a way as to provide us with Bjorken scaling at small $x$ (large $s$). The resulting amplitude (structure function) is applicable in the whole kinematical range, resonance region included. We call this unification “two dimensional duality” - one in $s$, the other one in $Q^2$.

At the early days of duality, off-mass continuation was attempted [8] by means of multi-leg (e.g. 6-point) dual amplitudes with ”extra” lines taken at their poles. Without going into details, here we only mention that scaling in this approach can be achieved [9] with nonlinear trajectories only, e.g. trajectories with logarithmic or constant asymptotics.

2 Notation and conventions

We use standard notation for the cross section and structure function (see Fig. 2):

$$\sigma^{\gamma^*p} = \frac{4\pi^2\alpha(1 + 4m^2x^2/Q^2)}{Q^2(1 - x)} \frac{F_2(x, Q^2)}{1 + R(x, Q^2)},$$

where $\alpha$ is the fine structure constant, $Q^2$ is the four-momentum transfer squared (with minus sign) or the momentum carried by the virtual photon, $x$ is the Bjorken variable and $s$ is the centre-of-mass energy squared of the $\gamma^*p$ system obeying the relation as follows:

$$s = Q^2(1 - x)/x + m_p^2,$$

where $m_p$ is the proton mass and $R(x, Q^2) = \sigma_L(x, Q^2)/\sigma_T(x, Q^2)$. For the sake of simplicity we set $R = 0$, that is a reasonable approximation.

We use the norm where

$$\sigma^{\gamma^*p}(s, t, Q^2) = \text{Im} A(s, t, Q^2).$$

According to the two-component duality picture [7], both the scattering amplitude $A$ and structure function $F_2$ are the sums of the diffractive and non-diffractive terms. At high energies both terms are of the Regge type. For $\gamma^*p$ scattering the positive-signature exchanges are allowed only. The dominant ones are the Pomeron and $f$ Reggeon, respectively. The relevant scattering amplitude is as follows (here $t = 0$):

$$A_i(s, Q^2) = i\beta_k(Q^2)\left(-\frac{s}{s_i}\right)^{\alpha_k(0)-1},$$

where $\alpha$ and $\beta$ are Regge trajectories and residues and $k$ stands either for the Pomeron or the Reggeon. As usual,
the residue is chosen to satisfy approximate Bjorken scaling for the structure function \[ F_2(x,Q^2) \sim (1-x)^n, \]
for high virtuality.

### 3 Nucleon resonances in inelastic electron-nucleon scattering

About thirty years ago Bloom and Gilman observed that the prominent resonances in inelastic electron-proton scattering do not disappear with increasing \( Q^2 \) with respect to the "background" but instead fall at roughly the same rate as any background. Furthermore, the smooth scaling limit proved to be an accurate average over resonance bumps seen at lower \( Q^2 \) and \( s \).

Since then, the phenomenon was studied in a number of papers and recently has been confirmed experimentally. These studies were aimed mainly to answer the questions: in which way a limited number of resonances can reproduce the smooth scaling behaviour?
According to the Veneziano (or resonance-reggeon) duality a proper sum of either t-channel or s-channel resonance exchanges accounts for the whole amplitude.

The main theoretical tools in these studies were finite energy sum rules and perturbative QCD calculations, whenever applicable. Our aim instead is the construction of an explicit dual model combining direct channel resonances, Regge behaviour typical for hadrons and scaling behaviour typical for the partonic picture.

The existence of resonances in the structure function at large $x$ close to $x = 1$ is not surprising by itself: as it follows from (1) and (2) they are the same as in $\gamma^* p$ total cross section, but in a different coordinate system.

The important question is whether and, if so, how a small number of resonances (or even a single one) can reproduce the smooth Bjorken scaling behaviour, known to be an asymptotic property which is typical for multiparticle processes.

The possibility that a limited (small) number of resonances can build up the smooth Regge behaviour was demonstrated by means of the finite-energy sum rules, but it was not realized explicitly like the Veneziano model (or its further modifications).

Actually, the early onset of Bjorken scaling, which is called "early, or precaution scaling", was observed with the first measurements of deep inelastic scattering at SLAC, where it was noticed that a more rapid approach to scaling can be achieved with the Bloom-Gilman (BG) variable $x' = x / (1 + m^2 x^2 / Q^2)$ instead of $x$ (or $\omega = 1 / x$). More recently the following generalization of the BG variable, such as

$$\xi = \frac{2x}{1 + \sqrt{1 + 4m^2x^2/Q^2}},$$

was suggested by O.Nachtmann [14]. We use the standard Bjorken variable $x$, however our results can be easily rewritten in terms of the above-mentioned modified variables.

First attempts to combine resonance (Regge) behaviour with Bjorken scaling were made [13,15,16] at low energies (large $x$), with the emphasis on the right choice of the $Q^2$-dependence, such as to satisfy the needed behaviour of form factors, vector meson dominance (VMD) with the requirement of Bjorken scaling. (N.B.: the validity (or failure) of the (generalized) VMD is still disputed). Similar
attempts in the high-energy (low \( x \)) region became popular recently, with the advent of the HERA data. They are presented in Sec. 5.

A consistent treatment of the problem requires the account for the spin dependence. For the sake of simplicity we ignore it in this paper (see e.g. [11]).

4 Factorization and dual properties

(bootstrap) of the vertices

Since the purpose of the present paper is the construction of a unified model realizing duality both in the \( s \) and \( t \) channels, we first attempt to identify its fragments, namely, the vertices (to be interpreted later on as \( Q^2 \)-dependent form factors).

Let us remind that the residue functions are completely arbitrary in the Regge pole model, but they are constrained in the dual model. We show this by using the low-energy (resonances) and high-energy (Regge) decomposition in the simple Veneziano model [3]

\[
V(s, t) = \int_0^1 dz z^{\alpha(s)}(1 - z)^{-\alpha(t)} = B(1 - \alpha(s), 1 - \alpha(t)) = \frac{\Gamma(1 - \alpha(s)) \Gamma(1 - \alpha(t))}{\Gamma(2 - \alpha(s) - \alpha(t))}.
\]

Furthermore,

\[
V(s, t) = \sum_{\alpha=0}^{\infty} \frac{1}{n - \alpha(s)} \frac{\Gamma(n + \alpha(t) + 1)}{n! \Gamma(\alpha(t) + 1)}.
\]

and, since for small \(|t|\) the \( \Gamma \) function varies slowly compared with the exponential one, the Regge asymptotic behaviour is

\[
V(s, t) \sim (\alpha' s)^{\alpha(t)},
\]

where \( \beta(t) = (\alpha')^{\alpha(t)} \) is the Regge residue.

Actually, one has to identify a single (and hence economic!) Regge exchange amplitude with a sum of direct-channel poles. Such an identification is not practical for infinite number of poles (e.g. the Veneziano amplitude) but, as we show below, is feasible if their number is finite (small). To anticipate the forthcoming discussion, we feed the \( Q^2 \)-dependence in the Regge residue at high energies (small \( x \)) and use the dual amplitude (with finite number of resonances!) for the whole kinematical region, including that of resonances. Relating the amplitude to the SF, we set \( t = 0 \).

To remedy the problems of the infinite number of narrow resonance, nonunitarity and lack of the imaginary part, we use a generalization of the Veneziano model free from the above-mentioned difficulties.

5 Dual amplitude with Mandelstam analyticity

The invariant dual on-shell scattering amplitude with Mandelstam analyticity (DAMA) applicable both to the diffractive and non-diffractive components reads [17]:

\[
D(s, t) = \int_0^1 dz \left( \frac{z}{g} \right)^{-\alpha(s' - 1)} \left( \frac{1 - z}{g} \right)^{-\alpha(t' - 1)} ,
\]

where \( s' = s(1 - z), \ t' = tz, \ g \) is a parameter, \( g > 1 \), and \( s, \ t \) are the Mandelstam variables.
For $s \to \infty$ and fixed $t$ it has the following Regge asymptotic behaviour

$$D(s, t) \approx \sqrt{\frac{2\pi}{\alpha_t(0)} g^{1+a+ib}} \left( \frac{s\alpha'(0) g \ln g}{\alpha_t(0)} \right)^{\alpha_t(0)-1},$$

(12)

where $a = Re \alpha\left(\frac{\alpha_t(0)}{\alpha'(0) \ln g}\right)$ and $b = Im \alpha\left(\frac{\alpha_t(0)}{\alpha'(0) \ln g}\right)$.

The pole structure of DAMA is similar to that of the Veneziano model except that multiple poles may appear at daughter levels. The presence of these multipoles does not contradict the theoretical postulates. On the other hand, they can be removed without any harm to the dual model by means the so-called Van der Corput neutralizer. The procedure [17] is to multiply the integrand of (11) by a function $\phi(x)$ with the properties:

$$\phi(0) = 0, \quad \phi(1) = 1, \quad \phi^n(1) = 0, \quad n = 1, 2, 3, ...$$

The function $\phi(x) = 1 - \exp \left( -\frac{x}{1-x} \right)$, for example, satisfies the above conditions and results [17] in a standard, "Veneziano-like" pole structure:

$$D(s, t) = \sum_n g^{n+\alpha_t(0)} \frac{C_n}{n - \alpha(s)},$$

(13)

where

$$C_n = \frac{\alpha_t(0) \left(\alpha_t(0) + 1\right) \ldots \left(\alpha_t(0) + n + 1\right)}{n!}.$$  

(14)

The pole term [13] is a generalization of the Breit-Wigner formula [1], comprising a whole sequence of resonances lying on a complex trajectory $\alpha(s)$. Such a "reggeized" resonances terminate.

Breit-Wigner formula has little practical use in the case of linear trajectories, resulting in an infinite sequence of poles, but it becomes a powerful tool if complex trajectories with a limited real part and hence a restricted number of resonances are used. Moreover, it appears that a small number of resonances is sufficient to saturate the direct channel.

Contrary to the Veneziano model, DAMA [11] not only allows but rather requires the use of nonlinear complex trajectories providing, in particular, for the imaginary part of the amplitude, resonance widths and resulting in a finite number of those. More specifically, the asymptotic rise of the trajectories in DAMA is limited by the condition (in accordance with an important upper bound derived earlier [18]):

$$\left| \frac{\alpha(s)}{\sqrt{s \ln s}} \right| \leq \text{const}, \quad s \to \infty.$$  

(15)

Actually, this upper bound can be even lowered up to a logarithm by requiring wide angle scaling behaviour for the amplitude.

The models of Regge trajectories combining the correct threshold and asymptotic behaviours have been widely discussed in the literature (see e.g. [19] for a recent treatment of this problem). A particularly simple model is based on a sum of square roots

$$\alpha(s) = \alpha_0 + \sum_i \gamma_i \sqrt{s_i - s},$$

where the lightest threshold (made of two pions or a pion and a nucleon) is important for the imaginary part, while the heaviest threshold limits the rise of the real part, where resonances terminate.

Dual amplitudes with Mandelstam analyticity with trajectories specified above are equally applicable to both: the diffractive and non-diffractive components of the amplitude, the difference being qualitative rather than quantitative. The utilization of a trajectory with a single thresh-
\[ \alpha_E(s) = \alpha_E(0) + \alpha_1 E(\sqrt{s_E} - \sqrt{s_E - s}) \] (16)

prevents the production of resonances on the physical sheet, although they are present on the nonphysical sheet, sustaining duality (i.e. their sum produces Regge asymptotic behaviour). This nontrivial property of DAMA makes it particularly attractive in describing the smooth background (dual to the Pomeron exchange) (see [17]).

The threshold value, slope and intercept of this exotic trajectory are free parameters.

For the resonance component a finite sum in (13) is adequate, but we use a simple model with the lowest threshold included explicitly and the higher ones approximated by a linear term:

\[ \alpha_R(s) = \alpha_R(0) + \alpha' s + \alpha_1 R(\sqrt{s_0} - \sqrt{s_0 - s}) \] (17)

where \( s_0 \) is the lowest threshold \( s_0 = (m_\pi + m_p)^2 \) in our case – while the remaining 3 parameters are adjusted to the known properties of relevant trajectories (\( N^* \) and \( \Delta \) isobar in our case). The termination of resonances, which are provided in DAMA by the limited real part, are effectively taken into account here by a cutoff in the summation of (13).

Finally, we note that a minimal model for the scattering amplitude is a sum

\[ A(s, t, u) = c(D(s, t) + D(u, t)) \] (18)

providing the correct signature at high-energy limit, \( c \) is a normalization factor. We disregard the symmetry (spin and isospin) properties of the problem, concentrating on its dynamics. In the limit \( s \to \infty, \ t = 0 \) we have \( u = -s \) and therefore

\[ A(s, 0, -s)|_{s \to \infty} = c \ D(s, 0)(1 + (-1)^{\alpha(0)-1}), \] (19)

where \( D(s, t) \) is given by eq. (12). For the total cross section in this limit we obtain:

\[ \sigma_T^\gamma p = \text{Im} \ A = C g^{\alpha(0)+a} (s \alpha'(0) \ln g)^{\alpha(0)-1}, \]

\[ \cdot (\sin(\alpha_t(0) - 1)\pi \cos(b \ln g) + (1 + \cos(\alpha_t(0) - 1)\pi) \sin(b \ln g)) \]

where \( C \) is a constant independent of \( s, \ g \) and \( \alpha'(0) \).

6 \( Q^2 \) dependence

Our main idea is to introduce the \( Q^2 \)-dependence in the dual model by matching its Regge asymptotic behaviour and pole structure to standard forms known from the literature. The point is that the correct identification of this \( Q^2 \)-dependence in a single asymptotic limit of the dual amplitude will extend it to the rest of the kinematical regions. We have two ways to do so, that is,

A) to combine Regge behaviour and Bjorken scaling limits of the structure functions (or \( Q^2 \)-dependent \( \gamma^* p \) cross sections), or

B) to introduce properly \( Q^2 \) dependence in the resonance region.

They should match to each other, if the procedure is correct, and the dual amplitude should take care of any further inter- or extrapolation.

It is obvious from eq. (4) that asymptotic Regge and scaling behaviour require the residue to fall like \( \sim (Q^2)^{-\alpha(0)+1} \).
Actually, it could be more involved if we require the correct $Q^2 \to 0$ limit to be respected and the observed scaling violation (the "HERA effect") to be included. Various models to cope with the above requirements have been suggested \[20,21,22\]. At HERA, especially at large $Q^2$, scaling is so badly violated that it may not be explicit anymore.

In combining Regge asymptotic behaviour with (approximate) Bjorken scaling, one can proceed basically in the following way – to keep explicitly a scaling factor $x^{-\Delta}$ (to be broken by some $Q^2$-dependence "properly" taken into account) \[21\].

\begin{equation}
F_2(x, Q^2) \sim x^{-\Delta(Q^2)} \left( \frac{Q^2}{Q^2 + Q_0^2} \right)^{1+\Delta(Q^2)},
\end{equation}

where $\Delta(Q^2) = \alpha_t(0) - 1$ may be a constant, in particular.

Note that since the Regge asymptotics of the Veneziano model is $\sim (-\alpha's)^{\alpha_t(0) - 1}$, the only way to incorporate there the $Q^2$-dependence is through the slope $\alpha'$. i.e. by making the trajectories $Q^2$-dependent, thus violating Regge factorization. The $Q^2$-dependent intercepts were used earlier \[21,22\] in a different context, namely, to cope with the observed "hardening" of the small-$x$ physics with increasing $Q^2$ (Bjorken scaling violation). Although we do not exclude this possibility (treating it as "effective" Regge pole), we study here the different option of introducing scaling violation in the constant $g$ appearing, besides $\alpha'$, in the residue of DAMA, eq (11).

Using the explicit Regge asymptotic form of DAMA, \[21\], and neglecting the logarithmic dependence of $g$, we make the following identification

\begin{equation}
g(Q^2)^{\alpha_t(0) + a} = \left( \frac{Q_{lim}^2}{Q^2 + Q_0^2} \right)^{\alpha_t(0)}.
\end{equation}

Note that eq. (22) is transcendent with respect to $g$, since $a = a(g) = Re \left( \frac{\alpha_t(0)}{\sigma(0) \ln g} \right)$. Another point to be mentioned is that this equation is not valid in the whole range of $Q^2$, since for $Q^2$ close to $Q_{lim}^2$, $g$ may get smaller than 1, which is unacceptable in DAMA. For large $Q^2$, the $Q^2$-dependence of $\log g$ and $b = b(Q^2) = Im \alpha \left( \frac{\alpha_t(0)}{\sigma(0) \ln g} \right)$ in eq. (20) cannot be neglected, it might contribute to scaling violation.

7 Scaling at large $x$

Let us now consider the extreme case of a single resonance contribution.

A resonance pole in DAMA contributes with

\begin{equation}
A(s, t) = g^{n + \alpha_t(0)} \frac{C_n}{n - \alpha(s)}.
\end{equation}
1.1 1.2 1.3 1.4 1.5

Fig. 5. $\gamma^* p$ total cross section as a function of $\sqrt{s}$. The dashed line shows the contribution from the $\Delta$ resonance, the dot-dashed line corresponds to the background, i.e. the contribution from the exotic trajectory. Here $Q^2 = 1 \text{ GeV}^2$.

At the resonance $s = s_{Res}$ one has $\text{Re} \alpha(s_{Res}) = n$ and $\frac{Q^2(1-x)}{x} = s_{Res} - m^2$, hence

$$F_2(x, Q^2) = \frac{Q^2(1-x)}{4\pi^2 \alpha \left(1 + \frac{4m^2x^2}{Q^2}\right)} \frac{C_n}{\text{Im} \alpha(s_{Res})} g(Q^2)^{n+\alpha(t)(0)}.$$  

As $x \rightarrow 1$, $Q^2 \approx \frac{s_{Res} - m^2}{1-x} \rightarrow \infty$ and

$$F_2(x, Q^2) \sim g \left(\frac{s_{Res} - m^2}{1-x}\right)^{n+\alpha(t)(0)}.$$  

By using the approximate solution $g(Q^2) \approx (Q^2_{lim}/Q^2)^{\alpha(t)(0)}$, where $a$ is a slowly varying function of $g$, we get for $x$ near 1

$$F_2(x, Q^2) \sim (1-x)^{\frac{\alpha(t)(0)x + \alpha(t)(0)}{\alpha(t)(0)+1}},$$  

where the limits for $x$ are defined by $Q^2_0 \ll \frac{s_{Res} - m^2}{1-x} \leq Q^2_{lim}.$

We recognize a typical large-$x$ scaling behaviour $(1-x)^N$ with the power $N$ (counting the quarks in the reaction) depending basically on the intercept of the $t$-channel trajectories.

8 Numerical Estimates

Having fixed the $Q^2$-dependence of the dual model by matching its Regge asymptotic behaviour to that of the structure functions, we now use this dual model to extrapolate it down to the resonance region, where its pole expansion (13) is appropriate - now complemented with the $Q^2$-dependence through $g(Q^2)$, fixed by eq. (22).

As has been already said, we write the imaginary part of the scattering amplitude as the sum of two terms - a diffractive (background) and non-diffractive (resonance) one. Note that $g(Q^2)$ has the same functional form (22) in both cases, only the values of the parameters differ (they are fixed from the small-$x$ fits (22) of the SF).

At low, resonance, energies the $\gamma^* p$ scattering exhibits a rich resonance structure intensively studied in a number...
of papers. About 20 resonances overlap, their relative importance varying with $Q^2$, but only a few can be identified more or less unambiguously. These are: $\Delta^+(1236)$ with $J^P = \frac{3}{2}^+$, $N^{*+}(1520)$, $J^P = \frac{3}{2}^+$, $N^{**+}(1688)$, $J^P = \frac{5}{2}^+$, and $N^{++}(1920)$ with $J^P = \frac{7}{2}^+$. They lie on the $\Delta$ and the exchange-degenerate $N$ trajectories. In this work we are mainly interested in introducing $Q^2$-dependence into the scattering amplitude, therefore we concentrate on a single resonance ($\Delta^+(1236)$) at different values of $Q^2$. We use trajectories in which the lowest pion-nucleon threshold is included explicitly, while higher thresholds are approximated by a linear term:

$$\alpha_{\Delta}(s) = 0.1 + 0.84s + 0.1331(\sqrt{s_0} - \sqrt{s_0 - s}),$$

where $s_0 = (m_\pi^2 + m_N^2)$. The above values of the parameters are chosen so as to fit the known mass and width of $\Delta^+(1236)$ in which the lowest pion-nucleon threshold is included explicitly, while higher thresholds are approximated by a linear term:

$$\alpha_{\Delta}(s) = 0.1 + 0.84s + 0.1331(\sqrt{s_0} - \sqrt{s_0 - s}),$$

where $s_0 = (m_\pi^2 + m_N^2)$. The above values of the parameters are chosen so as to fit the known mass and width of $\Delta^+(1236)$.
Table 1. Parameters used in the calculations shown in Figs. 4, 5 and 6.

|                  | $\Delta$ Resonance | Background |
|------------------|---------------------|------------|
| $Q_{\text{lim}}^2$, GeV$^2$ | 62                  | 120        |
| $Q_0^2$, GeV$^2$    | 0.01                | 2.5        |

Dual trajectory

- $\alpha_f(t)$ is dual to $\alpha_{\Delta}(s)$
- $\alpha_P(t)$ is dual to $\alpha_E(s)$

|                  | $\alpha_f(0) = 0.9$ | $\alpha_P(0) = 1 + 0.077 \cdot (1 + \frac{Q^2}{Q^2_{1.17}})$ |

Normalization coefficient $c = 0.03$.

the $\Delta$ resonance in a way consistent with the known linear parameterizations.

In the interval of interest $\sqrt{s} = 1.1 - 1.5$ GeV, $t = 0$, we have $u = m_X^2 - s < 0$, so it is far from resonance region, therefore we neglect the contribution from $D(u,t)$ for both the resonance and the background terms.

The smooth background is also modeled by a single term and exotic trajectory (16). As has been already explained, the direct channel Regge pole does not produce here physical resonances. The parameters of the exotic trajectory are:

$$\alpha_E(s) = -0.25 + 0.25(\sqrt{1.21} - \sqrt{1.21 - s}),$$ (23)

where $s_E = 1.1^2$ GeV$^2$ is an effective exotic threshold. Obviously "pole" does not mean a resonance in this case.

Figure 4 shows $g$ as a function of $Q^2$ for $\Delta$ and exotic trajectories. The resulting cross sections (imaginary part of the amplitude) in the resonance region is shown in Figs. 5 and 6 for two values of $Q^2 = 1$ and 6 GeV$^2$.

It is in qualitative agreement with the experimental data [4]. Figure 7 shows the dual properties of the cross section in two dimensions - one is the energy squared $s$ and the other one is virtuality $Q^2$. Table 1 shows the values of the parameters used in our calculations.

The main conclusions from our analysis are as follows:

A) the $Q^2$-dependence at low- and high-x (or high- and low-s) are interrelated and of the same origin;

B) even a single (low energy) resonance can produce a smooth scaling-like curve in the structure function (parton-hadron duality).

To summarize, we have suggested an explicit dual model in which the $Q^2$-dependence introduced in the low-$x$ domain is extended to the whole kinematic region, in particular, to the region of resonances. The resulting predictions for the first resonance in the $\gamma^*p$ system shown in Figs. 5, 6 are in quantitative agreement with data.

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