Ref.TH.2682-CERN

ROLE OF THE QUARK-PARTON MODEL IN DEUTERON INELASTIC SCREENING AT $t \neq 0$

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A B S T R A C T

We show that a naïve extension of the Gribov inelastic screening amplitude to $t \neq 0$ is affected by an inconsistency. This becomes apparent, if we represent the amplitude, as an integral of the six-point amplitude, and consists of a sharp discontinuity of the latter at the separation mass between the resonance region and the high-mass continuum. This discontinuity is caused by a damping factor appearing only in the high-mass continuum, described by a triple Regge formalism. We show that in a quark-parton model for hadronic diffraction, this factor is applied to the whole spectrum of intermediate diffractive states, not making any distinction between low and high masses. This prescription gives good agreement between theory and recent ISR data for pD elastic scattering at $|t| \sim 1 \text{ (GeV/c)}^2$.

*) Supported in part by INFN.

1. INTRODUCTION

In a recent analysis of proton-deuteron and deuteron-deuteron elastic scattering data from the CERN Intersecting Storage Rings (ISR)¹⁾, evidence was found for inelastic screening at momentum transfers $|t| \sim 0.1-0.5 \text{ (GeV/c)}^2$. As for the case of the total cross-section defect²⁾, the inelastic states are described by a low-mass series of resonant states and by a high-mass continuum. However, a naïve extension at $t \neq 0$ of the inelastic screening amplitude for $t = 0$, leads to an anomalous discontinuity of the corresponding six-point amplitude^{3, 4)}. This results from the presence of a form factor in the hadron-Pomeron coupling, which damps the high-mass part at high t , while the low-mass part remains unchanged. This discontinuity is responsible for a strong violation of the duality sum rule⁵⁾, the amount of violation being a quantity rising with t .

Continuity and duality are restored if even the low-mass part is damped by the same form factor. The presence of this form factor for the whole spectrum of intermediate diffractive states is required by a quark-parton model for diffractive reactions⁶⁾, in the first order, and it is interpreted as a Fourier transform of the transverse probability distribution of the interacting partons inside the hadron.

This also provides a nice interpretation in terms of the quark-parton model of the hadron-Pomeron vertex function, clearly related to the compositeness of the hadron.

The plan of the paper is the following: in Section 2 we state the problem in the assumption of Pomeron dominance, and we compare data with theory, with and without the form factor applied in the low-mass region. In Section 3 we calculate in a quark-parton model the t -distribution for the reaction $pp \rightarrow p + X$ and the inelastic screening amplitude. From the comparison between the two, we find that the inelastic screening amplitude contains the form factor of the parton distribution.

2. THE PROBLEM

The formula for the inelastic screening amplitude in the case of hadron-deuteron elastic scattering is

$$\begin{aligned} f_{IS}(q_T) = & \frac{i}{2\pi p_{Lab}} \int d^2 q'_T \sum_n f_n\left(\frac{\vec{q}_T}{2} - \vec{q}'_T\right) f_n\left(\frac{\vec{q}_T}{2} + \vec{q}'_T\right) S(\vec{q}'_T, q_L) \\ & + \frac{i}{2\pi p_{Lab}} \int d^2 q'_T dM_X^2 \Delta F_6^{3P}\left[q_T^2, \left(\frac{\vec{q}_T}{2} - \vec{q}'_T\right)^2, \left(\frac{\vec{q}_T}{2} + \vec{q}'_T\right)^2\right] S(\vec{q}'_T, q_L), \end{aligned} \quad (2.1)$$

where p_{lab} is the laboratory momentum of the incident hadron; f_n the resonance production amplitudes; q_T, q_L the transverse and longitudinal part of the momentum transfer; M_X the mass of the intermediate state; and ΔF_6^{3P} the discontinuity of the six-point amplitude³⁾, in the hypothesis of Pomeron dominance. For the purpose of giving a qualitative discussion, it is not necessary to introduce other Regge poles; this complication would make our treatment more obscure, without qualitatively changing our results.

As usual we may neglect in the integrals of (2.1) the $\vec{q}'_T$ dependence in the production amplitudes, and we may write the sum over n as an integral over M_X^2 . The result is a sum of two integrals in M_X^2 :

$$\begin{aligned} f_{IS} = & -\frac{i p_{Lab}}{2\pi^2} \int_{M_0^2}^{\bar{M}^2} dM_X^2 \frac{d\sigma}{dt dM_X^2} \int d^2 q'_T S(\vec{q}'_T, q_L) \\ & + \frac{i}{2\pi p_{Lab}} \int_{\bar{M}^2}^{\infty} dM_X^2 \Delta F_6^{3P}\left(q_T^2, \frac{q_T^2}{4}, \frac{q_T^2}{4}\right) \int d^2 q'_T S(\vec{q}'_T, q_L) \end{aligned} \quad (2.2)$$

where M_0 is the threshold mass and $\bar{M}$ the separation mass between the low-mass resonant spectrum and the continuum. In the low-mass integral we have used the relation

$$\frac{d\sigma}{dt dM_\pi^2} = -\frac{\pi}{p_{Lab}^2} f_\pi^2$$

which is valid in the assumption of purely imaginary amplitude⁷⁾. As is well known, $t = -q_T^2$ at small angles. The above formula cannot be reduced to a compact form, similar to the Gribov⁷⁾ expression for the total cross-section defect, since the relation between ΔF_6 and $d\sigma/dtdM_X^2$ is

$$\Delta F_6^{3P}(q_T^2, \frac{q_T^2}{4}, \frac{q_T^2}{4}) = -e^{-\alpha q_T^2} \frac{p_{Lab}^2}{\pi} \frac{d\sigma}{dt dM_X^2} \quad (2.3)$$

where the factor $\exp(-\alpha q_T^2)$ comes from the Pomeron residuum, which depends on the external momenta (Fig. 1). In the following we call it R , to recall that it is connected to the residuum and to keep open a possibility of having a parametrization, different from the usual Gaussian. This factor can be determined, for instance, from the energy-independent part of the t -dependence in proton-proton elastic scattering or from the energy-independent part of the t -dependence of $d\sigma/dtdM_X^2$, in the assumption that the triple-Pomeron vertex does not depend on t .

If we join the two integrals in (2.2) and consider them as a unique integral of one M_X^2 function, this function is affected by a strong discontinuity at the separation mass $\bar{M}$, as shown by Fig. 2^{*}). Following the Russian literature³⁾, this function is identified with the absorptive part of the six-point amplitude and it should therefore be continuous in M_X^2 . One might think that this discontinuity is apparent, because we have implicitly taken into account the integral on the complex cut; we have shown, however, that this contribution is small and may be neglected altogether⁴⁾.

^{*}) In the figure we use for the low-mass region a parametrization with two BW forms⁸⁾ and for the high-mass region the usual $1/M^2$ form.

This discontinuity also causes a violation of the duality sum rule⁵⁾, which should be valid also for the inelastic screening case.

In order to restore all these good properties in the framework of only Pomeron exchange, we may apply the R factor also to the low-mass integral of (2.2). This recipe is giving a fairly good agreement with ISR data¹⁾ at $|t| \sim 1 \text{ (GeV/c)}^2$, as shown in Fig. 3 where we plot $d\sigma/dt$, calculated using the extended form of Glauber theory [Eq. (2.2)] normalized to the ordinary Glauber prediction, and compare it with experimental data normalized in the same way. Neglecting the R factor in the low-mass integral we get the dashed line, clearly off the data at $|t| \sim 1 \text{ (GeV/c)}^2$. Neglecting R in both integrals of (2.2) would restore all the good properties we discussed above, but the results shown as the dot-dashed line in Fig. 3 are in very bad agreement with the data.

As we shall see in the following section, the recipe is easily justified in the quark-parton model.

3. INELASTIC SCREENING IN A QUARK-PARTON MODEL

It was repeatedly shown in the literature that elastic and inelastic diffraction can be reproduced by a nuclear model for the hadron, where the constituents are sometimes identified as quarks⁹⁾ or gluons¹⁰⁾ or more generally as partons. Recently Miettinen and Pumplin⁶⁾ succeeded in describing both proton elastic and inelastic diffraction, representing the proton as a superposition of free parton states, eigenstates of the T-matrix. Their model describes, with only slight modifications, the inelastic screening in the deuteron. As we shall see, the factor R comes naturally in the model at the first order and is interpreted as a Fourier transform of the probability distribution of the partons.

In order to study the R factor, we have first of all to study the inclusive distribution for the reaction $pp \rightarrow p + X$. As described at length in Ref. 6, this is obtained as dispersion of the T matrix

$$\frac{d\sigma}{dt} = \frac{1}{4\pi} \left[\langle \tilde{t}^2 \rangle - \langle \tilde{t} \rangle^2 \right],$$

where

$$\begin{aligned} \langle \tilde{t}^j \rangle = & \sum_{N=0}^{\infty} e^{-G^2} \frac{G^{2N}}{N!} \int \prod_{i=1}^N d^2 b_i dy_i |C(\vec{b}_i, y_i)|^2 \times \\ & \left(\int d^2 b e^{i\vec{q}_T \cdot \vec{b}} \left\{ 1 - \prod_{k=1}^N [1 - \tau(\vec{b}_k + \vec{b}, y_k)] \right\} \right). \end{aligned} \quad (3.1)$$

Here the notation is exactly the same as in Ref. 6: for completeness we recall that G^2 is the average number of partons, $\vec{b}_i$ are the transverse coordinates of the partons with respect to the central coordinate $\vec{b}$ of the proton, and y_i are their rapidities; $|C(\vec{b}_i, y_i)|^2$ is the probability distribution of the parton, and the τ 's are the scattering functions of the partons with the potential describing the target proton. It is important to notice that in the model of Ref. 6 the number of constituents is not fixed, and this makes the diffractive cross-section (3.1) behave as an exponential at small q_T^2 .

After summing over N , Miettinen and Pumplin⁶⁾ obtain the following expression for the dispersion (3.1):

$$\begin{aligned} \frac{d\sigma}{dt} = & \frac{1}{4\pi} \int d^2 b d^2 b' e^{i\vec{q}_T \cdot (\vec{b} - \vec{b}')} \exp \left[-G^2 (\langle \tau(\vec{b}) \rangle + \langle \tau(\vec{b}') \rangle) \right] \times \\ & \left[\exp(G^2 \langle \tau(\vec{b}) \tau(\vec{b}') \rangle) - 1 \right], \end{aligned} \quad (3.2)$$

where

$$\langle \tau(\vec{b}) \tau(\vec{b}') \rangle = \int d^2 b_1 dy_1 |C(\vec{b}_1, y_1)|^2 \tau(\vec{b}_1 + \vec{b}, y_1) \tau(\vec{b}_1 + \vec{b}', y_1).$$

A similar expression holds for $\langle \tau(\vec{b}) \rangle$.

In this formalism one may easily obtain the expression of the inelastic screening amplitude in proton-deuteron elastic scattering:

$$f_{IS} = -\frac{i p_{Lab}}{2\pi} \langle \psi | \int d^2b e^{i\vec{q}_T \cdot \vec{b}} \exp \left\{ -G^2 \left[\langle \tau(\vec{b} + \frac{\vec{s}}{2}) \rangle + \langle \tau(\vec{b} - \frac{\vec{s}}{2}) \rangle \right] \right\} \\ \times \left\{ \exp \left[G^2 \langle \tau(\vec{b} + \frac{\vec{s}}{2}) \tau(\vec{b} - \frac{\vec{s}}{2}) \rangle \right] - 1 \right\} | \psi \rangle, \quad (3.3)$$

where $\vec{s}$ is now the transverse position of the relative coordinate of the two nucleons, which act as two distinct potentials, one at position $\frac{1}{2} \vec{s}$ and the other at $-\frac{1}{2} \vec{s}$; ψ is the deuteron ground-state wave function.

The derivation of (3.3) is given in detail in the Appendix: here we give some comments to explain the essential differences between formulas (3.2) and (3.3). The most important difference is in the phase factor: in (3.3) the factor $\exp(i\vec{q}_T \cdot \vec{b})$ results from the product of the two phase factors $\exp[i(\vec{q}_T/2)(\vec{b} + \vec{s}/2)]$ and $\exp[i(\vec{q}_T/2)(\vec{b} - \vec{s}/2)]$, related to the two scatterings, each one having momentum transfers $\vec{q}_T/2$. In this case the two phases have the same sign, while in (3.2) the corresponding phases have different signs: the reason is simply that while (3.3) is an amplitude, (3.2) is a cross-section, and the second elementary amplitude is complex conjugated.

To study the question of the R factor in detail, we consider both (3.2) and (3.3) in the first order in G^2 ; the result for (3.2) is

$$\frac{d\sigma}{dt} = \frac{1}{4\pi} G^2 \ll \int d^2b d^2b' e^{i\vec{q}_T \cdot (\vec{b} - \vec{b}')} \tau(\vec{b} + \vec{b}_1) \tau(\vec{b}' + \vec{b}_1) \gg_{(\vec{b}_1, y_1)} \quad (3.4)$$

where we indicate with the subscript $(\vec{b}_1, y_1)$ that the average is made on both $\vec{b}_1$ and y_1 . The two average symbols indicate that the average is made on two variables.

Substituting the Fourier transforms of the τ functions^{*)}, we get

$$\begin{aligned}
 \frac{d\sigma}{dt} &= \frac{G^2}{4\pi(2\pi)^4} \ll \int d^2b d^2b' e^{i\vec{q}_T \cdot (\vec{b} - \vec{b}')} \int d^2q'_T e^{-i\vec{q}'_T \cdot (\vec{b} + \vec{b}_1)} \times \\
 &\quad \int d^2q''_T e^{-i\vec{q}''_T \cdot (\vec{b}' + \vec{b}_1)} \tilde{\tau}(\vec{q}'_T) \tilde{\tau}(\vec{q}''_T) \gg_{(\vec{b}_1, y_1)} \\
 &= \frac{G^2}{4\pi} \ll e^{-i\vec{q}_T \cdot \vec{b}_1} e^{i\vec{q}_T \cdot \vec{b}_1} \tilde{\tau}^2(\vec{q}_T) \gg_{(\vec{b}_1, y_1)} \quad (3.5) \\
 &= \frac{G^2}{4\pi} \langle \tilde{\tau}^2(\vec{q}_T) \rangle_{(y_1)}.
 \end{aligned}$$

Here only the average on y_1 is necessary, because $\tilde{\tau}$ does not depend on $\vec{b}_1$.

Doing the same for (3.3), we get

$$\begin{aligned}
 p_{IS} &= -\frac{iP_{lab} G^2}{2\pi} \langle \psi | \ll \int d^2b e^{i\vec{q}_T \cdot \vec{b}} \tau(\vec{b} + \vec{b}_1 + \frac{\vec{s}}{2}) \tau(\vec{b} + \vec{b}_1 - \frac{\vec{s}}{2}) \gg_{(\vec{b}_1, y_1)} | \psi \rangle \\
 &= -\frac{iP_{lab} G^2}{2\pi} \frac{1}{(2\pi)^4} \langle \psi | \ll \int d^2b e^{i\vec{q}_T \cdot \vec{b}} \int d^2q'_T e^{-i\vec{q}'_T \cdot (\vec{b} + \vec{b}_1 + \frac{\vec{s}}{2})} \tilde{\tau}(\vec{q}'_T) \times \\
 &\quad \int d^2q''_T e^{-i\vec{q}''_T \cdot (\vec{b} + \vec{b}_1 - \frac{\vec{s}}{2})} \tilde{\tau}(\vec{q}''_T) \gg_{(\vec{b}_1, y_1)} | \psi \rangle \quad (3.6) \\
 &= -\frac{iP_{lab} G^2}{2\pi} \frac{1}{(2\pi)^2} \langle \psi | \ll e^{-i\vec{q}_T \cdot \vec{b}_1} \int d^2q'_T e^{i\vec{q}'_T \cdot \vec{s}} \tilde{\tau}(\frac{\vec{q}_T}{2} - \vec{q}'_T) \tilde{\tau}(\frac{\vec{q}_T}{2} + \vec{q}'_T) \gg_{(\vec{b}_1, y_1)} | \psi \rangle.
 \end{aligned}$$

*) We follow Ref. 6 in the normalization of the Fourier transform

$$\tau(\vec{b}) = \frac{1}{(2\pi)^2} \int d^2q'_T e^{-i\vec{q}'_T \cdot \vec{b}} \tilde{\tau}(\vec{q}'_T).$$

Making the usual assumption¹¹⁾ of neglecting the momentum dependence of the scattering functions $\tilde{f}$, with respect to the form factor $\langle \psi | \exp(i\vec{q}_T' \cdot \vec{s}) | \psi \rangle$, we get, using (3.5)

$$\begin{aligned} p_{IS} &= - \frac{i p_{Lab}}{2\pi} G^2 \frac{1}{4\pi} \left\langle \frac{1}{\tau^2} \right\rangle \left\langle \tilde{\tau}^2 \left(\frac{\vec{q}_T}{2} \right) \right\rangle_{(y_1)} \left\langle e^{-i\vec{q}_T \cdot \vec{b}_1} \right\rangle_{(\vec{b}_1)} \\ &= - \frac{i p_{Lab}}{2\pi} \left\langle \frac{1}{\tau^2} \right\rangle \left. \frac{d\sigma}{dt} \right|_{\frac{t}{4}} \left\langle e^{-i\vec{q}_T \cdot \vec{b}_1} \right\rangle_{(\vec{b}_1)}. \end{aligned} \quad (3.7)$$

The comparison between this expression and formula (2.3) leads to the identification

$$R = \left\langle e^{-i\vec{q}_T \cdot \vec{b}_1} \right\rangle_{(\vec{b}_1)},$$

confirming the validity of the hypothesis of Section 2, where this factor was also applied to the low-mass integral in (2.2). Indeed, we sum here on all possible intermediate diffractive states, not making any distinction between low and high masses. It is instructive to notice that if we neglect the dependence of τ on b_1 , we do not get the damping factor in (3.6): in other words the existence of the form factor is due to the internal structure of the incident hadron.

In the model of Ref. 6, the parton probability distribution is Gaussian and therefore

$$\left\langle e^{-i\vec{q}_T \cdot \vec{b}_1} \right\rangle_{(\vec{b}_1)} = e^{-\beta q_T^2/4},$$

with $\beta = 6 \text{ (GeV/c)}^{-2}$. The determination of this parameter, using the energy-independent part of the measured slope for proton-proton scattering, leads to $\alpha \sim 2 \text{ (GeV/c)}^{-2}$, comparable to the number determined from Ref. 6, $\beta/4 = 1.5 \text{ (GeV/c)}^{-2}$.

If we extract this number from the inclusive cross-section, and subtract the energy-dependent part exactly, we get $\alpha = 1.5 \text{ (GeV/c)}^{-2}$.

4. CROSS-SECTIONS OF DIFFRACTIVE STATES

The conclusions we have reached in the elastic case could be extended to the inelastic one, in particular to the intriguing problem of coherent production¹²⁾. There we have intermediate inelastic states which may have low or high mass exactly as in the elastic case. For the high-mass integral, there is a damping factor due to the overlap between the initial and final parton distributions. If the same damping factor is applied to the low-mass integral, it modifies the interaction cross-section of the final diffractive states, which is measured by evaluating the size of the low-mass integral consistent with experiment. So the actual interaction cross-section of diffractive states should be obtained multiplying the measured one by the inverse of the damping factor,

$$\sigma_T(N^*N) = \sigma_T(\text{measured}) e^{\alpha' \langle q_T^2 \rangle},$$

where $\langle q_T^2 \rangle$ is the average value of the momentum transfer in the particular experiment.

For the recent ISR experiment with deuteron beams¹³⁾ the average value of $|t|$ that is sensitive to re-scattering is $\langle q_T^2 \rangle = 0.32 \text{ GeV}^2$. If $\alpha' = \alpha = 2 \text{ GeV}^{-2}$, then the factor $\exp(\alpha' \langle q_T^2 \rangle) = 1.92$, so that the actual value of the $\sigma_T(N^*N)$ would be around the more reasonable value of 40 mb.

5. CONCLUSIONS

We have shown in this paper that in a quark parton model the inelastic screening amplitude contains a damping factor which comes from the confinement of the interacting partons in the transverse coordinates. This confinement occurs in all the s-channel absorptive models for diffractive reactions, but it is also common to more evolutionary dynamical descriptions¹⁴⁻¹⁶⁾, using the concept of colour, and it is based on the observation of a transverse momentum cut-off in particle production.

This damping factor already exists in the usual triple Regge formulation of the inelastic screening amplitude, but does not appear in the low-mass resonant part of the spectrum, causing unpleasant discontinuity in the six-point amplitude and violation of the duality sum rule. The presence of this factor has phenomenological implications for the calculations of inelastic screening at $t \neq 0$ and could partially explain the low value of the effective cross-section in coherent production experiment.

Acknowledgements

G.A. is grateful to S. Pokorski for many illuminating discussions and to T. Ericson for a critical reading of the manuscript. F.B. is grateful to the CISRIF, Trieste, for financial support.

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APPENDIX

In this appendix we calculate formulas (3.2) and (3.3) in detail. The cross-section for inelastic diffraction (3.2) is defined as the standard deviation of the t-matrix, given by the fluctuations of its eigenvalues in the diffractive sector:

$$\begin{aligned} \frac{d\sigma}{dt} &= \frac{1}{4\pi} \left[\langle |\tilde{t}(\vec{q}_T)|^2 \rangle - |\langle \tilde{t}(\vec{q}_T) \rangle|^2 \right] \\ &= \frac{1}{4\pi} \int d^2b \int d^2b' e^{i\vec{q}_T \cdot (\vec{b} - \vec{b}')} \left\{ \langle t(\vec{b}) t^*(\vec{b}') \rangle - \langle t(\vec{b}) \rangle \langle t^*(\vec{b}') \rangle \right\} \end{aligned} \quad (A.1)$$

Here we have used the definition

$$\tilde{t}(\vec{q}_T) = \int d^2b e^{i\vec{q}_T \cdot \vec{b}} t(\vec{b})$$

given in Ref. 6.

The derivation is very similar in the inelastic screening case, but for the presence of the deuteron wave function $\psi(\vec{r})$. The two nucleons in the deuteron are in the positions $\vec{b}' = \vec{b} + (\vec{s}/2)$ and $\vec{b}'' = \vec{b} - (\vec{s}/2)$, which means that the difference $\vec{b}' - \vec{b}''$ is constrained to be equal to the relative transverse coordinate $\vec{s}$ of the nucleons in the deuteron:

$$\begin{aligned} f_{IS} &= -\frac{iP_{Lab}}{2\pi} \langle \psi | \int d^2b' \int d^2b'' \delta(\vec{b}' - \vec{b}'' - \vec{s}) e^{i\frac{\vec{q}_T}{2} \cdot (\vec{b}' + \vec{b}'')} \\ &\quad \left\{ \langle t(\vec{b}') t(\vec{b}'') \rangle - \langle t(\vec{b}') \rangle \langle t(\vec{b}'') \rangle \right\} | \psi \rangle \\ &= -\frac{iP_{Lab}}{2\pi} \langle \psi | \int d^2b e^{i\vec{q}_T \cdot \vec{b}} \left\{ \langle t(\vec{b} + \frac{\vec{s}}{2}) t(\vec{b} - \frac{\vec{s}}{2}) \rangle - \langle t(\vec{b} + \frac{\vec{s}}{2}) \rangle \langle t(\vec{b} - \frac{\vec{s}}{2}) \rangle \right\} | \psi \rangle \end{aligned} \quad (A.2)$$

We assume here that the total transverse momentum $\vec{q}_T$ is equally shared between the two scatterers: this assumption gives the right form for the Glauber elastic amplitude, as is clear from the subtracted term in (A.2).

Figure captions

- Fig. 1 : Triple Regge diagram for the deuteron inelastic screening. The mass of the upper Reggeon is $-q_T^2$, contrary to the case of inclusive production $pp \rightarrow p + X$, where $q_T^2 = 0$.
- Fig. 2 : Function of M_X^2 appearing in the M_X^2 integral of inelastic screening amplitude (2.2) at $q_T^2 = 0.65 \text{ (GeV/c)}^2$. The separation mass is taken to be $\bar{M} = 2 \text{ GeV}$: the parametrization of the low-mass part is from Ref. 8. α is taken equal to half of the energy-independent slope of $(d\sigma/dtdM^2)$ ($pp \rightarrow p + X$) at FNAL energies.
- Fig. 3 : The ratio $(d\sigma/dt)/(d\sigma/dt)_{\text{Glauber}}$, compared with ISR data of Ref. 1 at the highest energy ($\sqrt{s} = 63 \text{ GeV}$), normalized in the same way. The dot-dashed line shows the result, neglecting the R factor altogether. The dashed line corresponds to the recipe (2.2) and (2.3); that is, with the R factor only in the high-mass part. The continuous line is the result with R applied also to the low-mass part.

Fig. 1

Fig. 2

Fig. 3