INSTITUTE OF THEORETICAL AND EXPERIMENTAL PHYSICS

4

ITEP-1 – 93

C.Bourrely, V.L.Morgunov, O.V.Kancheli,
K.A.Ter-Martirosyan

SUPERCRITICAL POMERON, ELASTIC SCATTERING AND PARTICLES PRODUCTION AT SUPERHIGH ENERGIES

MOSCOW 1993

УДК 530.I45 М-I6

SUPERCRITICAL POMERON, ELASTIC SCATTERING AND PARTICLES PRODUCTION AT SUPERHIGH ENERGIES: Preprint ITEP 93-1/

C.Bourrely[*)], V.L.Morgunov, O.V.Kancheli, K.A.Ter-Martirosyan - 8p.

New approach to the rescatterings on the supercritical Pomeron is developed. Rescattering vertices γ_n are determined in the form corresponding to the contribution of two diffraction particles beams.

Their general properties are discussed and different particles production cross sections are written in terms of γ_n.

Some examples of the simultaneous description of high energy elastic scattering, diffraction production and pionization processes are considered.

Fig. - 5, ref. - 4

Centre de Physique Théorique, CNRS-Luminy, 13288 Marseille Cedex 9, France

1. The general rescattering picture

At extremely high energies $s/m_N^2 \to \infty$ the supercritical Pomeron theory (having $\Delta = \alpha_P(0) - 1 > 0, \Delta \simeq 0.1 - 0.2$) determines the elastic scattering amplitude $M(s,t) = \int_0^\infty \frac{i}{2} F(b,s) \cdot J_o(bp_\perp) b db$ in the form of the sum of a large number ($< n > \sim (s/m_N^2)^\Delta$) of rescatterings: $F(b,s) = 1 - S(b,s)$,

$$S(b,s) = exp[2i\delta(b,s)] = \sum_{n=0}^{\infty} \frac{(-v(b,s))^n}{n!} \gamma_n(b,\xi) \tag{1}$$

where γ_n are some unknown rescattering Gribov's coefficients;

$$v(b,s) = \frac{2}{i} \int M_P(p_\perp, s) J_o(p_\perp b) p_\perp dp_\perp$$

-is the one Pomeron contribution, $M_P \simeq i\lambda_P(p_\perp^2) exp[\xi\Delta - \alpha'_P(0)\xi_1 p_\perp^2]$, $\xi_1 = ln\, s/m_N^2 - i\pi/2$ and the decreasing dependence of the Pomeron residue $\lambda_P(p_\perp^2) = g^{(1)}(p_\perp^2) g^{(2)}(p_\perp^2)$ on $t = -p_\perp^2$ is not determined by the theory. The function $S(b,s)$ will be called below "S - matrix". We remind that the Gaussian form of the residue $\lambda_P = \lambda_P^o exp(-R^2 p_\perp^2)$, valid only for small $p_\perp^2 < M^2$, leads to a simple $v(b,s) = (\lambda_P^o/r_o^2) exp(\xi\Delta - b^2/4r_o^2)$, $r_o^2 = R_P^2 + \alpha'_P \xi$, which is at $\xi \gg 1$ very large for small $b < b_o$ (where $v \sim exp(\xi_1 \Delta) \simeq (s/m_N^2)^\Delta \gg 1$) and is small: $v \ll 1$, at large b. [1]

Considering γ_n as an analytic function of n : $\gamma_n = \gamma(n)$ at $n = 1, 2, \ldots$ and demanding that $S = S(v)$ in Eq.(1) has to decrease for $v \to \infty$, as this corresponds to the black disk absorption $F(b,s) \to 1$, at small $b < b_c$ and at $s/m_N^2 \to \infty$ (see Fig.1 where $b_o = b_o(\xi) \approx a_o \xi$ is the interaction radius) it is easy to show that $\gamma(n)$ can have singularities only at the left half-plane of the complex n-plane [2] and, the conditions must hold :

$$\gamma(0) = 1, \quad \gamma(1) = 1 . \tag{2}$$

[1] At $b^2/4r_o^2 = \xi\Delta$, i.e. at $b = b_o$, $b_o^2 = 4(R_P^2 + \alpha'_p \xi)\xi\Delta$, the value of $v(b,s)$ is of order of unity. Thus, the interaction radius $b_o = 2\sqrt{\alpha'_P \Delta}\xi$ rise like $ln(s/m_N^2)$, what leads to the Froissart limit: $\sigma^{tot} \simeq 2\pi b_o^2 \simeq (8\pi\alpha'_p\Delta)\xi^2$ in this model.

[2] Authors thank M.Braun for this remark. General features of $\gamma(n)$ will be considered in a separate publication.

Note that:

i) The dependence of γ_n on $\xi \simeq ln(s/m_N^2)$ is not essential at all, as this dependence can be reproduced by a small change of the parameter Δ in $v \sim exp(\xi\Delta) \sim (s/m_N^2)^\Delta$.

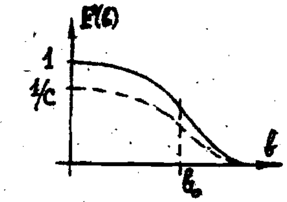

Fig.1 *The profile F(b) of the scattering, dashed line shows the result of quasieikonal approach.*

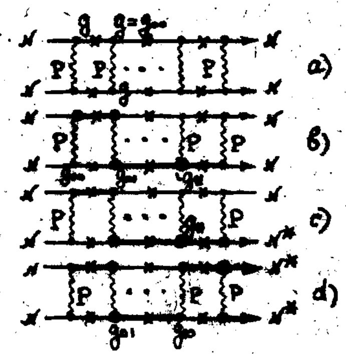

Fig.2 *Eikonal model graphs a), b), c), d)-two chanal picture.*

ii) The so called eikonal approximation $\gamma_n \equiv 1, n \geq 1$, corresponding to the graphs shown in Fig.2a, leads to zero values of diffraction dissociation (DD) cross section $\sigma^{DD} = 2\sigma_{1D} + \sigma_{2D}$ and gives $S = exp[-v(b,s)], F = 1 - exp(-v)$. However, even in this approximation, the Harward-Marseille group [1] was able to obtain a perfect description of pp, $\bar{p}p$ elastic scattering up to $p_\perp^2 \sim 2(\frac{Gev}{c})^2$ choosing the power type residue

$$\lambda_P(p_\perp^2) = \lambda_P^o \cdot [(1 + p_\perp^2/m_1^2)(1 + p_\perp^2/m_2^2)]^{-2} \frac{1 - p_\perp^2/æ^2}{1 + p_\perp^2/æ^2}, \tag{3}$$

where the first $p_\perp^2$ depending factor corresponds to a matter distribution in the proton analogous to a charge distribution, with $m_1 \simeq 0.58 GeV, m_2 \simeq 1.7 GeV$ and the zero at large $p_\perp = æ$ in $\lambda_P(p_\perp^2)$ was introduced in [1] by hand ($\lambda_P^o \simeq 2.1 (GeV/c)^{-2}$, $æ \simeq 2.0\ GeV/c$).

iii) The "quasi-eikonal" approximation [3], which corresponds to $\gamma_n = C^{n-1}$, $C = 1 + \sigma^{DD}/\sigma^{el}$, $n \neq 0$ and to $\gamma_o = 1 - C^{-1}$, i.e. to $F = \frac{1}{C}(1 - exp(-Cv))$, was successfully used [2] to describe, simultaneously, the

$pp, \bar{p}p$ scattering at small $p_\perp^2 \leq p_o^2 \leq 0.2 - 0.3(Gev/c)^2$ and, also, all available information about the particle production at high energies. However it does not describe the $d\sigma^{el}/dp_\perp^2$ behaviour at $p_\perp^2 \geq p_o^2$ and instead of the black disk at $s/m_N^2 \to \infty, v \to \infty$ (at $b < b_o$), leads to the constant grey disk (see dashed line in Fig.1) due to violation of the condition $\gamma(0) = 1$.

Below, all particle production cross-section will be expressed in terms of the coefficients $\gamma_n(b)$ and some simple models will be considered, where $\gamma_n \neq 1$ arises due to the production, in the rescattering process, of two diffraction beams of particles with small masses $M_1 \sim M_2 \sim m_N$ (see Fig.3a). The particular model of such a beam in the $NN - N\bar{N}$ interaction cases can be the production of an excited state N^* of the nucleon in rescatterings (as is shown in Fig.2b). Fig.2c and 2d show graphs describing one jet and two jets DD processes.

2. Particles production cross-sections

Our normalization corresponds to $\sigma^{tot} = 8\pi Im M(s,0)$, $d\sigma^{el}/dp_\perp^2 = 4\pi \mid M(s,p_\perp^2) \mid^2$. Writing any type of cross section in the form

$$\sigma^i(s) = \int \tilde{\sigma}^i(b,s) d^2b = 2\pi \int_0^\infty \tilde{\sigma}^i(b,s) b db$$

one finds: $\tilde{\sigma}^{tot}(b) = 2[1 - ReS(v)]$, $\tilde{\sigma}^{in}(b) = 1 - \mid S(v) \mid^2$ and also $\tilde{\sigma}^{el}(b) = \mid F(b) \mid^2 = \mid 1 - S(v) \mid^2$. Now, using AGK rules one obtains, (with $v = v_1 + iv_2$), the cross section $\tilde{\sigma}_k(b)$ for the production of some number $k = 1, 2, \ldots$ of Pomeron showers of hadrons in the form:

$$\tilde{\sigma}_k(b) = \frac{(2v_1)^k}{k!} \sum_{n=0}^\infty \frac{(-2v_1)^n}{n!} \gamma_{n+k}, \quad \tilde{\sigma}^{pion}(b) = \sum_{n=1}^\infty \tilde{\sigma}_n(b) = 1 - S(2v_1) \qquad (4)$$

where $S(2v_1)$ is just the sum Eq.(1) with $2v_1$ substituted for $v(b,\xi)$. The difference

$$\tilde{\sigma}^{DD}(b) = \tilde{\sigma}^{in}(b) - \tilde{\sigma}^{pion}(b) = S(2v_1) - \mid S(v) \mid^2 \qquad (5)$$

is the DD cross-section $\tilde{\sigma}^{DD} = 2\tilde{\sigma}_{1D} + \tilde{\sigma}_{2D}$ - the sum of the cross-sections of one and two sides diffraction beams or jets production. In S - unitary theory $\tilde{\sigma}^{DD}(b)$ and all other $\tilde{\sigma}^i(b)$ naturally must be positive for all values of b.

Considering the square of the modulus of the amplitude corresponding to the graphs shown in Fig.3b and 3c one obtains DD cross-sections $\tilde{\sigma}_{1D}(b)$ and

$\tilde{\sigma}_{2D}(b)$ for NN scattering in the form of the following sums:

$$\tilde{\sigma}_{1D} = \sum_{n_1,n_2=1}^\infty \sigma_{n_1,n_2}[\sqrt{\gamma_{n_1+n_2}\gamma_{n_1}\gamma_{n_2}} - \gamma_{n_1}\gamma_{n_2}],$$

$$\tilde{\sigma}_{2D} = \sum_{n_1,n_2=1}^\infty \sigma_{n_1,n_2}[\sqrt{\gamma_{n_1+n_2}} - \sqrt{\gamma_{n_1}\gamma_{n_2}}]^2 \qquad (6)$$

where $\sigma_{n_1,n_2} = \frac{(-v)^{n_1}(-v^*)^{n_2}}{n_1!n_2!}$. Note that the term $2\sqrt{\gamma_{n_1+n_2}\gamma_{n_1}\gamma_{n_2}}$ cancels out in the sum $2\tilde{\sigma}_{1D} + \tilde{\sigma}_{2D}$ and Eq.(5) is thus reopened.

3. Particles beams at rescatterings and γ_n factors; some attempts of data fit

There exist many functions $\gamma_n(b) = \gamma(n,b)$ of variable n which satisfy Eqs.(2) and have singularities only in the left half n-plane. To have some orientation in the form of $\gamma_n = \gamma(n)$ let us consider the physical picture of production of two diffraction beams of particles with small effective mass $M_1 \sim M_2 \sim m_N$ in the rescattering process shown in Fig.3: for elastic scattering (Fig.3a), for one DD jet production (Fig.3b) and for two DD jets production (Fig.3c).

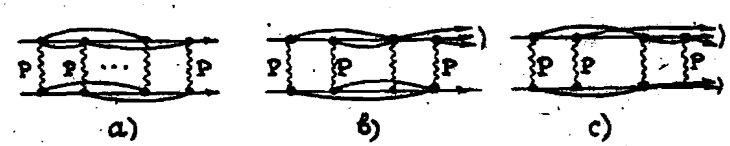

Fig.3 *Two difraction beams pictures:*
a) for elastic scattering, *b) for* σ_{1D}, *c) for* σ_{2D}.

The upper and lower beams are separated by a large rapidity interval; we can consider their effect separately putting $\gamma_n = \beta_{1n}\beta_{2n}$, or $\gamma_n = \beta_n^2$ for NN interaction case. The longer is the length of these beams (i.e. the more links $k \leq n - 1$ they pass in apper or in lower parts in Fig.3) the smaller must be [3]

[3]As the smaller is the probability for the beam to reach the finite targer with area $\pi R_0^2 (R_0 \sim 1/m_N)$ in the diffusion in the impact parameter space.

their contribution $\sim \frac{\lambda}{k+a}$ to the eikonal value $\beta_n = 1$, where $\lambda, a \leq 1$ are some parameters. Therefore:

$$\beta_n = \beta(n) = \prod_{k=1}^{n-1}(1+\frac{\lambda}{k+a})^{n-k} \quad , \quad with \ \ \beta_0(0) = \beta(1) = 1 \ , \qquad (7)$$

as, for $k = 1$ the enhanced factor $C_1 = 1 + \frac{\lambda}{1+a}$, which coresponds to each link in Fig.3, enters $n-1$ times. Comparing this factor with old "quasieikonal" model [2] where $\gamma_2 \simeq C_1^2 \simeq 1 + 1/2$ one concludes that $\lambda/(1+a) \simeq 1/4$. The factor $C_2 = 1 + \lambda/(2+a)$ enters $n-2$ times in β_n and corresponds to beams avoiding one Pomeron vertex, factor $C_3 = 1 + \lambda/(3+a)$ corresponds to beams avoiding two vertices in Fig.3 and enters $n-3$ times e.t.c. This results in Eq.(7).

The function $\beta(n)$ in Eq.(7) can be shown to have the simple pole at $n = -n_0, n_0 = 1 + a$ and has at $n \gg 1$ the asymptotic form

$$\beta(n) = [\Gamma(n+n_0)/\Gamma(n_0)n_0^n]^\nu \ , \quad \nu/n_0 \simeq 1/4 \qquad (8)$$

(or, simply, $\beta = (\Gamma(n+1))^\nu, \nu \sim 1/4$ at $a = 0, n_0 = 1$). For $\gamma_n = \beta^2(n)$ in Eq.(1) it leads to $S(v) \sim (lnv)/v^{n_0}$ at $v(b,\xi\ ) \gg 1$ what seems quite reasonable physically. However, the b-dependence of γ_n factors is lost at all in Eqs.(7), (8) (note, by the way, that equalities $\beta(0) = \beta(1) = 1$ still holds in the form (8)).

To reproduce this dependence let us remind the old two channel (N and N^*) picture of Fig.2b,c,d resulting [3] in $\gamma_n = \gamma'(n) = (\beta'(n))^2$ with

$$\beta'(n) = \frac{\varepsilon\eta^n + 1}{1+\varepsilon}\left(\frac{1+\varepsilon}{1+\eta\varepsilon}\right)^n = t_1^n\lambda_1 + t_2^n\lambda_2 \qquad (9)$$

where $\varepsilon, \eta \sim 1$ are some DD-parameters determining (see [3] for details) the form of 2×2 - Pomeron vertex $\hat{g}_P$, and (9) is obtained by trivial diagonalization of it: $\hat{g}_P \rightarrow (\hat{g}_P^{diag})_{ij} = g_i(p_\perp^2)\delta_{ij}, g_i = t_i g_0(p_\perp^2), i,j = 1,2$, using diagonal "beam" states $\varphi_i = a_1^{(i)}N + a_2^{(i)}N^*$. In Eq.(9) $\lambda_1 = \varepsilon\lambda_2 = \varepsilon/(1+\varepsilon), t_1 = \eta t_2, t_2 = (1+\varepsilon)/(1+\varepsilon\eta)$ and $\lambda_1 + \lambda_2 = 1, \lambda_1 t_1 + \lambda_2 t_2 = 1$. Considering each of beams in Fig.3 as consisting of an infinite number of similar diagonal states φ_i one obtains the more general form:

$$\beta'(n) = \sum_i \lambda_i \cdot t_i^n \qquad (10)$$

where λ_i and t_i are some parameters constrained by the conditions $\sum_i \lambda_i = \sum_i \lambda_i t_i = 1$. Putting in Eq.(1) $\gamma_n' = (\beta'(n))^2$ one obtaines :

$$S(b,s) = \sum_{i,j} \lambda_i\lambda_j exp[-t_i t_j v_{ij}(b,s)] \qquad (11)$$

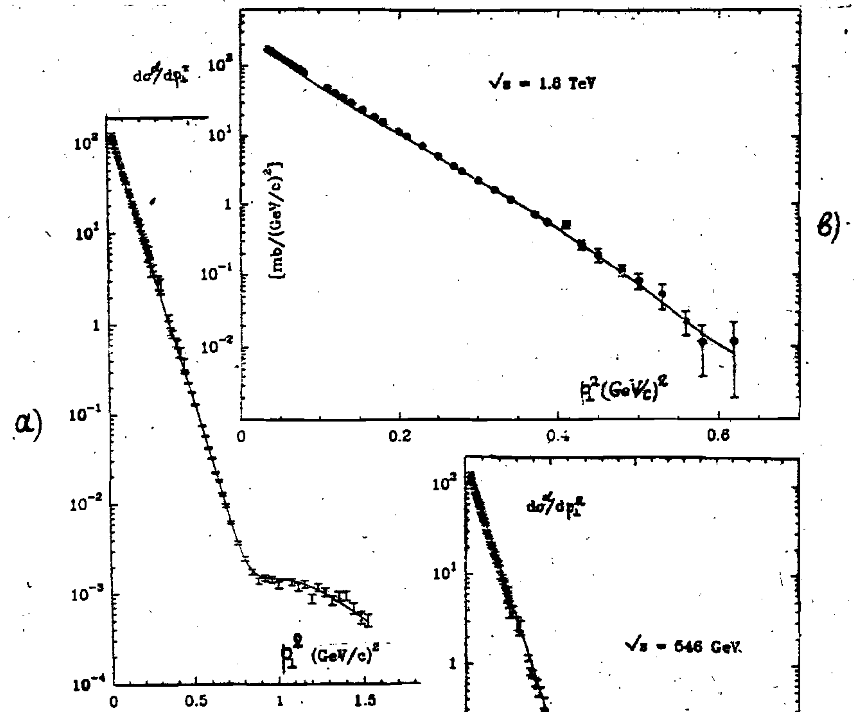

Fig.4 $d\sigma^{el}/dp_\perp^2$ *behaviour for* $S(v) = (1 + v/n_0)^{-n_0}$*: a) for* $\sqrt{s} = 546\ GeV$*; for* $\sqrt{s} = 1.8 TeV$ *- б).*

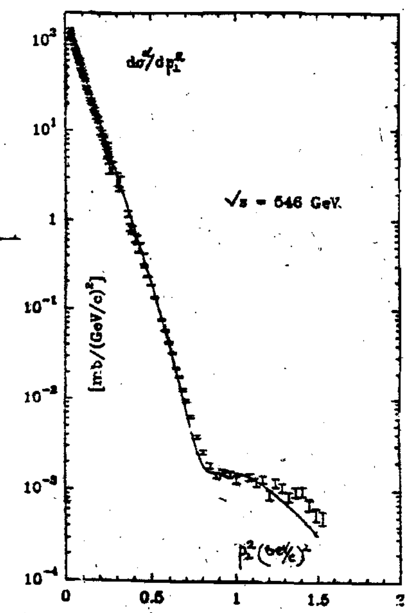

Fig.5 *Fit of* $d\sigma^{el}/dp_\perp^2$ *for* $\gamma_n = (\beta'(n))^2\gamma_n^{(1)}$.

where it was taken into accounts that Pomeron vertices $g_i(p_\perp^2)$ for different diagonal beam states can have different dependence on $p_\perp^2$ (i.e. can have different parameters m_1, m_2 in Eq.(3) ; $\lambda_P(p_\perp^2) \rightarrow \lambda_P^{ij} \sim g_i g_j$ in v_{ij}). This will just reproduce effectively the b-dependence of γ_n- coefficients.

Note, that Eq.(10) is the particular case of a more general representation:

$$\beta'(n) = \int_0^\infty \lambda(t) t^n dt \qquad \text{for} \quad \lambda(t) = \sum_i \lambda_i \delta(t - t_i) \ .$$

We note also, that the simplest form for $\gamma_n = \beta^2(n)$ is given by Eq.(8) at $\nu = 1/2$: $\gamma_n = \gamma_n^{(1)} = \frac{\Gamma(n+n_0)}{\Gamma(n_0)n_0^n}$; in Eq.(1) it leads to the simple S-matrix $S(v) = (1 + v/n_0)^{-n_0}$ discussed by Braun and Pajares [4] (with, $-n$ at the place of n.)

Eqs.(7)-(11) can be used for describtion of $d\sigma^{el}/dp_\perp^2$ and of particle production data at high energy. Two simplest examples are considered below:

a) Using $S(v) = (1 + v/n_0)^{-n_0}$ one obtains the following set of the best values of parameters: $n_0 = 2.27$, while in the Pomeron amplitude: $\Delta = 0.1$, and $\alpha'_P(0) = 0.03, \lambda_P^0 = 2.9$ in $(GeV/c)^{-2}$ and also in Eq.(3) $m_1 = 0.58, m_2 = 13.7, \ae = 16.6$ in GeV/c. The $d\sigma^{el}/dp_\perp^2$ behaviour is very good and is shown in Fig.4a for $\sqrt{s} = 546 GeV$ and in Fig.4b for $\sqrt{s} = 1.8 TeV$. The resulting values of σ^{tot} are: 63 and 74 mb corresrondingly, while values of $\rho = (ReM/ImM)_{p_\perp=0}$ are: 0.106 and 0.100.

However, too small σ^{DD} value was obtained here from Eq.(5): at $\sqrt{s} = 546\ GeV$, $\sigma^{DD} = 2\sigma_{1D} + \sigma_{2D} \simeq 2.4 mb$ was find as compared with $(\sigma^{DD})_{exp} \simeq 0.42\sigma^{el} \simeq 5.5mb$.

b) Using $\gamma_n = (\beta'(n))^2 \gamma_n^{(1)}$ with the same $\gamma_n^{(1)} = \Gamma(n + n_0)/\Gamma(n_0)n_0^n$ as above and $\beta'_0(n)$ in the form (9) one finds the parameters (in the same units as above): $n_0 = 3.92, \Delta = 0.112, \alpha'_P = 0.03, \lambda_P^0 = 4.66; m_1 = 0.75, m_2 = 13.7, \ae = 16.6$ - being taken the same in both poles $t_1 = 2.92, t_2 = 0.55$ (what corresponds to $\epsilon = 4.31, \eta = 0.19$, or to $\lambda_1 = 0.19, \lambda_2 = 0.81$ in Eq.(9)). The fit of $d\sigma^{el}/dp_\perp^2$ at $\sqrt{s} = 546 GeV$, shown in Fig.5, is not so good as in Fig.4, however $\sigma^{DD} = 4.1mb$ is here much closer to the experiment $(\sigma^{DD}(546))_{exp} \simeq 5.5mb$.

These first result are promising and show that calculations have to be extendet to a more realistic version of γ_n, e.g. defined in Eqs.(9), (10) with different parameters m_1, m_2 of Pomeron vertices for different $v_{ij}(b, s)$ in Eq.(11).

Authors thank K.Boreskov, M.Braun, A.Kaidalov, K.Pajares, J.Soffer for discussions and useful remarks.

References

[1] C.Bourrely, J.Soffer, T.T.Wu//Phys.Rev. D19 3249 (1979); Nucl.Phys. B247, 15 (1984); Z.Phys. C37, 369 (1988)

[2] A.B.Kaidalov, K.A.Ter-Martirosyan//Phys.Lett. 117B, 247 (1982); Sov.Nucl.Phys. 39, 1545 (1984); 40, 211 (1984)

[3] K.A.Ter-Martirosyan//Nucl.Phys. A477, 696 (1988)

[4] M.Braun, K.Pajares//Phys. Lett. (in press) Santjago preprint US FT 91/B/1991.

К.Бурели и др.

Сверхкритический померон, упругое рассеяние и рождение частиц при сверхвысоких энергиях.

Подписано к печати 21.12.92 Формат 60х90 1/8 Офсетн.печ.
Усл.-печ.л.0,5. Уч.-изд.л.0,3. Тираж 260 экз. Заказ I
Индекс 3649

Отпечатано в ИТЭФ, 117259, Москва, Б.Черемушкинская, 25