acta phys. slov. 36 (1986) No. 4

A NOTE ON TRANSITIONS BETWEEN QUARKS AND GLUONS[1]

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Using effective Lagrangian models we derive some new and interesting low-energy theorems for matrix elements of scalar gluonic current between physical (quarkonium) states of the pseudoscalar meson nonet.

ЗАМЕЧАНИЕ О ПЕРЕХОДАХ МЕЖДУ КВАРКАМИ И ГЛЮОНАМИ

На основе модельных эффективных лагранжианов в работе выведены некоторые новые интересные низкоэнергетические теоремы для матричных элементов скалярных глюонных токов между физическими (кваркониевыми) состояниями нонета псевдоскалярных мезонов.

I. INTRODUCTION

The knowledge of matrix elements for transitions between ordinary (i. e. quarkonium) hadronic states caused by pure gluonic currents (for a review, see, e. g. [1]) is interesting not only theoretically but also from an experimental point of view, e. g. for a clear identification of some experimentally found particles as gluonic bound states (i. e. glueballs or gluonia). Within perturbative QCD transitions betweeen gluonic and quark degrees of freedom are suppressed by the factors $O(\alpha_s)$, $\alpha_s = g^2/4\pi$, g being the strong coupling constant. These factors are responsible for the so-called OZI rule [2] formulated even before the QCD era. It forbids the quark line annihilation and is phenomenologically supported, for example, by the smallness of the $\Phi \to \pi\pi$ decay, by an approximate equality of the ϱ and the ω meson masses, etc.

A very popular and simple theoretical formulation of this phenomenon is within the multicolour chromodynamics [3]. In this approach the number of colours N_c is very large and instead of the expansion in α_s the expansion in N_c^{-1} is used, $N_c\alpha_s$ being constant for $N_c \to \infty$. Here, the OZI suppression is related to positive powers in $1/N_c$, e. g. the transition between gluonium and two quark mesons is suppressed as $O(1/N_c)$. Thus, as applied to gluonium decay, this theory generally predicts [4] that glueball states must be very narrow with widths of tens MeV. However, it has been argued [1] that due to the dominance of nonperturbative effects there are exceptions from the OZI rule, especially in the scalar and pseudoscalar channels. In the 0^- channel, for instance, large radiative decays [5] of the J/Ψ meson into $\gamma\eta'$ and $\gamma\eta$ (proceeding through an intermediate gluonic stage [6]) are examples of such strong couplings between $\eta'(\eta)$ and gluons. In the 0^+ channel an exception from the OZI rule can be represented by the dominant decay of meson $\Psi(3685)$ into $\pi\pi\, J/\Psi$ [5] with final pions produced by the S-wave. This decay is shown [7] to proceed in two stages: first $\Psi(3685) \to ggJ/\Psi$ and then two gluons gg are converted into $\pi\pi$. In this way the decay has been satisfactorily explained [7] by using the following low-energy theorem [7]

$$\langle \pi^+(p_1)\pi^-(p_2)|H(O)|O\rangle|_{\text{chiral limit}} = -q^2 + O(q^4), \tag{1}$$

where q^2 is the invariant $(\text{mass})^2$ of the $\pi^+\pi^-$ system and $H(x)$ is the scalar gluonic current to be specified later. Thus, in agreement with experiment eq. (1) explicitly demonstrates the disappearance of any suppression factors (like $O(\alpha_s)$) in coupling between gluons and quarks in the o^+ channel.

The purpose of this paper is to rederive eq. (1) in a different manner by using effective chiral Lagrangians (for a review and further references see, e. g. [8–9, 17]). In this way we shall also find new and interesting generalizations of eq. (1) for the case when the matrix element on the l. h. s. of eq. (1) contains any physical pseudoscalar states (i. e. with real masses) from the whole nonet (section II.). Some conclusions are given in section III.

II. MATRIX ELEMENTS OF THE SCALAR GLUONIC CURRENT BETWEEN PSEUDOSCALAR MESON STATES

The trace of the energy-momentum tensor of QCD has been proved to be of the form [10]

$$(\Theta^{\mu}_{\mu})_{\text{QCD}} = -\frac{b}{8}\frac{\alpha_s}{\pi} F^{(a)}_{\mu\nu} F^{(a)\mu\nu} + \sum_i m_i \bar{q}_i q_i, \tag{2}$$

where $F^{(a)}_{\mu\nu}$'s are gluon field strength tensors, $b = 11 - 2N_F/3$, $N_F = 3$ is the number of light quark flavours, m_i and $q_i(x)$ being the mass and field, respectively, of a quark of flavour i, ($i = u, d, s$). Here, we have neglected contributions due to the anomalous dimension of $\bar{q}q$ operators as well as $O(\alpha_s^2)$ contributions to the first term in eq. (2). Since we are interested in low energy physics of pseudoscalar

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particles, heavy quark flavours will also be neglected. The trace of the energy-momentum tensor of QCD is a renormalization group-invariant quantity, and we easily see that in a chiral symmetry limit (i. e., if $m_i=0$) it is directly related to the scalar gluonic current $H(x)$:

$$H(x)=\frac{9}{8}\frac{\alpha_s}{\pi}F^2(x)\equiv\frac{9}{8}\frac{\alpha_s}{\pi}F^{(a)}_{\mu\nu}F^{(a)\mu\nu}(x), \tag{3}$$

where $b=9$. Thus, to find low-energy theorems of type (1), we just need to obtain analogical relations for matrix elements of the trace Θ^μ_μ between pseudoscalar meson states. In the interesting low-energy region this can be done very effectively by using phenomenological Lagrangians [8, 9, 17].

i) Nonlinear effective Lagrangians

We shall start our considerations assuming that the low-energy dynamics of the nonet of pseudoscalar mesons is described by the following generally accepted nonlinear phenomenological Lagrangian [9, 17].

$$\mathcal{L}_{NL}=\frac{1}{4}\mathrm{Tr}\,[(\partial_\mu U)(\partial^\mu U^+)]+\frac{m_0^2f^2}{48}[\mathrm{Tr}(\ln U-\ln U^+)]^2-\frac{1}{4}\mathrm{Tr}\,[M(U+U^+)], \tag{4}$$

where M is proportional to the 3×3 quark mass matrix, the pion decay constant $f=93$ MeV, and m_0 is related to the masses of pseudoscalar mesons as follows

$$m_0^2=m_{\eta'}^2+m_\eta^2-2m_k^2 \tag{5}$$

$U(x)$ is parametrized as the unitary matrix:

$$U(x)=f\exp\left(\mathrm{i}\sum_{j=0}^{8}\frac{\lambda_j\varphi_j(x)}{f}\right), \tag{6}$$

where φ_i' s ($i=0, 1, \ldots, 8$) are fields for the nonet of pseudoscalar mesons, and Gell-Mann λ matrices are normalized to $\mathrm{Tr}\,(\lambda_i\lambda_j)=2\delta_{ij}$. The second term in eq. (4) is required by the axial anomaly and breaks explicitly the axial U (1) symmetry but conserves the chiral SU (3) × SU (3) one thus giving the mass $m_0\neq0$ to the pseudoscalar η' particle even in the chiral limit (i. e., when $M=0$). In this way the U (1) problem is solved [9, 17].

Rewriting the Lagrangian (4) in terms of the scalar u_j' s and pseudoscalar v_j' s ($j=0, 1, \ldots, 8$) fields defined by relations

$$u_j=\frac{1}{4}\mathrm{Tr}\,[\lambda_j(U+U^+)], \tag{7}$$

$$v_j=\frac{1}{4\mathrm{i}}\mathrm{Tr}\,[\lambda_j(U-U^+)]$$

and assuming that the fields u and v have dimensions equal to a real number d (i. e., under dilatation transformations $x\to\varrho x$, $u\to\varrho^{-d}u$ and $v\to\varrho^{-d}v$), one obtains an „improved" [11] energy-momentum tensor from eq. (4) as follows [12]

$$\Theta_{\mu\nu}=\sum_{i=0}^{8}[(\partial_\mu u_i)(\partial_\nu u_i)+(\partial_\mu v_i)(\partial_\nu v_i)]-g_{\mu\nu}\mathcal{L}_{NL}+ \\ +\frac{d}{6}[g_{\mu\nu}\Box-\partial_\mu\partial_\nu]\sum_{i=0}^{8}(u_i^2+v_i^2). \tag{8}$$

We see that this tensor differs from the canonical one by the last term in eq. (8) that is required for getting a simple connection between the trace of eq. (8) and divergence of the dilatation current $\mathcal{D}_\mu(x)$ [11], i. e.,

$$\partial_\mu\mathcal{D}^\mu=\Theta^\mu_\mu \tag{9}$$

It is worth noting that the dilatation "charge" $D=\int d^3x\,\mathcal{D}^0(x)$ generates dilatation transformations, e. g., for field u with dimension d we have

$$[D(t),u(x)]_{t=x^0}=-\mathrm{i}(x^\mu\partial_\mu+d)u(x). \tag{10}$$

However, in the present special case of parametrization (6) the fields u and v have dimension $d=0$ [12–14] and the difference between eq. (8) and the canonical form of the energy-momentum tensor disappears. Then, with operators in the normal order and due to eqs. (4), (6) and (7), eq. (8) becomes

$$(\Theta_{\mu\nu})_4=:\left\{\sum_{i=0}^{8}[(\partial_\mu u_i)(\partial_\nu u_i)+(\partial_\mu v_i)(\partial_\nu v_i)]- \right. \\ \left. -g_{\mu\nu}\mathcal{L}_{NL}\right\}:+\frac{1}{4}g_{\mu\nu}\langle\Theta^\lambda_\lambda\rangle_{0/\text{chiral limit}}, \tag{11}$$

where index "4" labels the correspondence to eq. (4) and the constant term is added in eq. (11) to achieve a correct normalization of $\langle 0|\Theta^\mu_\mu|0\rangle$ in the chiral limit. With the equations of motion, the trace of eq. (11) is as follows

$$(\Theta^\mu_\mu)_4=:-\frac{1}{2}\mathrm{Tr}\,[(\partial_\mu U)(\partial^\mu U^+)]-\frac{m_0^2f^2}{12}[\mathrm{Tr}(\ln U-\ln U^+)]^2+ \\ +\mathrm{Tr}[M(U+U^+)]:+\langle\Theta^\mu_\mu\rangle_{0/\text{chiral limit}}. \tag{12}$$

We easily see that VEV of eq. (12) leads to a trivial identity in the chiral-symmetry limit thus justifying the normalization chosen in eq. (11). Now from eqs. (2), (3) and (12) and after some manipulations we obtain the following effective relation for the gluonic current (3):

$$H = :\frac{1}{2}\,\mathrm{Tr}[(\partial_\mu U)(\partial^\mu U^+)] + \frac{m_0^2 f^2}{12}[\mathrm{Tr}\,(\ln U - \ln U^+)]^2 - $$
$$-\frac{3}{4}\,\mathrm{Tr}\,[M(U+U^+)]: + \frac{9}{8}\left\langle \frac{\alpha_s}{\pi} F^2 \right\rangle_{0/\mathrm{chiral\ limit}}. \tag{13}$$

From eqs. (3) and (13) we get

$$\left\langle \frac{\alpha_s}{\pi} F^2 \right\rangle_0 = \left\langle \frac{\alpha_s}{\pi} F^2 \right\rangle_{0/\mathrm{chiral\ limit}} + \frac{8}{3} f^2 \left(m_K^2 + \frac{1}{2} m_\pi^2 \right), \tag{14}$$

where VEV of the mass term was estimated from eqs. (4) and (6). Eq. (14) is very interesting and valuable because it gives us an idea how the gluon condensate is changed if one proceeds from the chiral symmetry limit to the real world and vice versa. For instance, for the value [15] $\langle (\alpha_s/\pi) F^2 \rangle_0 = 0.012\ \mathrm{GeV}^4$ one obtains $\langle (\alpha_s/\pi)\, F^2 \rangle_{0/\mathrm{chiral\ limit}} = 0.006\ \mathrm{GeV}^4$, and we notice a large difference between these two values. In other words, a chirally symmetrical world does not seem to be a good approximation in the case of the gluon condensate. However, when the phenomenological estimation of $\langle (\alpha_s/\pi) F^2 \rangle_0$ has to be larger by a factor $2 \div 3$ [16], then the chiral limit is still a reasonable approximation of the real world.

Due to eqs. (4 – 7) and (14) there are no free parameters in eqs. (11 – 13), and we can calculate any matrix element of these operators between pseudoscalar meson states. As usual, calculations have to be done in the tree approximation, and we shall also use the covariant normalization of states:

$$\langle p | p' \rangle = (2\pi)^3 2\omega_p \delta^{(3)}(p - p'). \tag{15}$$

In this way from eqs. (11) and (12) we get, for example, the following relations

$$\langle P(p_1) \bar{P}(p_2) | \Theta_{\mu\nu} | 0 \rangle = \frac{1}{2}(r_\mu r_\nu - q_\mu q_\nu + g_{\mu\nu} q^2), \tag{16}$$

and

$$\langle P(p_1) \bar{P}(p_2) | \Theta_\mu^\mu | 0 \rangle = q^2 + 2m_P^2, \tag{17}$$

where $r = p_1 - p_2$, $q = p_1 + p_2$, m_P is the mass of particle P and $P\bar{P} = \pi^+\pi^-$, K^+K^-, etc. Analogously, from eq. (13) one obtains the new theorems as follows

$$\langle \eta(p_1)\eta(p_2) | (-H(0)) | 0 \rangle = q^2 + m_\eta^2 + (m_{\eta'}^2 + m_\eta^2 - 2m_K^2)\sin^2 \Phi$$
$$\langle \eta'(p_1)\eta(p_2) | H(0) | 0 \rangle = (m_{\eta'}^2 + m_\eta^2 - 2m_K^2)\cos\Phi \sin\Phi, \tag{18}$$
$$\langle \eta'(p_1)\eta'(p_2) | (-H(0)) | 0 \rangle = q^2 + m_{\eta'}^2 + (m_{\eta'}^2 + m_\eta^2 - 2m_K^2)\cos^2\Phi,$$

where the $\eta\eta'$ mixing angle Φ is given by [9, 17]

$$\mathrm{tg}\, 2\Phi = -\frac{4\sqrt{2}(m_K^2 - m_\pi^2)}{3[m_0^2 - 2(m_K^2 - m_\pi^2)/3]} \tag{19}$$

and m_0 is from eq. (5). This leads to the value $\Phi = -18°$ in a good agreement with the most recent independent theoretical [18] and experimental [19] determinations.

Thus, we see that eq. (12) (or (13)) effectively (in an operator form) represents a generalization of eq. (1) not only for nonzero masses but also for the whole nonet of pseudoscalar mesons. In fact, from eq. (12) (or 13)) any interesting matrix element of the operator Θ_μ^μ (or H) between physical pseudoscalar states can be calculated in a straightforward and easy way. Some of the results of such calculations are explicitly given by eqs. (16 – 18).

ii) Linear effective Lagrangians

While within the nonlinear phenomenological Lagrangian [9, 17] approach the validity of eq. (1) is directly obvious from eqs. (6) and (13), the situation is not so clear in the case of linear effective Lagrangian models (for a review and further references, see, e. g., [8]). We shall show this starting with the simple linear σ – model containing three pseudoscalar pion $\pi_i(x)$ ($i = 1, 2, 3$) and one scalar $\sigma(x)$ fields. They form the (1/2, 1/2) representation of the chiral $SU(2) \times SU(2)$ group and the chirally symmetric Lagrangian is given as follows [8]

$$\mathcal{L}_L = \frac{1}{2}(\partial_\mu \sigma)^2 + \frac{1}{2}\sum_{i=1}^{3}(\partial_\mu \pi_i)^2 - \frac{1}{2}\mu^2\left(\sigma^2 + \sum_{i=1}^{3}\pi_i^2\right) -$$
$$-\lambda\left(\sigma^2 + \sum_{i=1}^{3}\pi_i^2\right)^2. \tag{20}$$

where λ and μ^2 are suitable parameters. With a conventional assignment of dimension 1 to $\sigma(x)$ and $\pi_i(x)$ and using equations of motion we obtain

$$(\Theta_\nu^\nu)_{20} = \mu^2\left(\sigma^2 + \sum_{i=1}^{3}\pi_i^2\right) \tag{21}$$

from the most general form of $\Theta_{\mu\nu}$ (eq. (8)). Spontaneous breaking of chiral symmetry through the existence of $\langle 0 | \sigma | 0 \rangle = \sigma_0 \neq 0$ leads to the necessity to correct the theory by the following redefinition of field σ:

$$\sigma(x) = \sigma_0 + \sigma'(x). \tag{22}$$

Correcting eq. (20) in this way and eliminating the term linear in $\sigma'(x)$ from it by using the vacuum stability condition ($\mu^2 = -4\lambda\sigma_0^2$) we arrive at the right Lagrangian. This Lagrangian gives zero masses for pions, a nonzero mass for σ-particle ($m_\sigma \neq 0$) and the following interaction term

$$\mathscr{L}_{\sigma\pi\pi}(x) = -\frac{m_\sigma^2}{2\sigma_0}\,\sigma'(x)\sum_{i=1}^{3}\pi_i^2(x). \tag{23}$$

In terms of the $\sigma'(x)$ field eq. (21) reads

$$(\Theta_\nu^\nu)_{20} = \mu^2\sigma_0 - m_\sigma^2\sigma_0\sigma' - \frac{m_\sigma^2}{2}\left(\sigma'^2 + \sum_{i=1}^{3}\pi_i^2\right). \tag{24}$$

It seems clear from eqs. (23) and (24) that dominant contributions to the matrix element $\langle \pi^+(p_1)\pi^-(p_2)|\Theta_\nu^\nu|0\rangle$ (which are obviously proportional to m_σ^2 instead of to $q^2=(p_1+p_2)^2$) break the validity of eq. (1) [14]. However, more precise calculations show that such dominant contributions cancel each other leading again to eq. (1). In fact, we have

$$\langle \pi^+(p_1)\pi^-(p_2)|(\Theta_\nu^\nu)_{20}|0\rangle = -m_\sigma^2 - \frac{m_\sigma^4}{q^2-m_\sigma^2} = q^2 + 0(q^4) \tag{25}$$

in agrement with eq. (1). Very recently eq. (25) has also been proved [20] in another way, by using the relation

$$(\Theta_\nu^\nu)_{20} = \mu^2\left(\sigma^2 + \sum_{i=1}^{3}\pi_i^2\right) - 4\lambda\left(\sigma^2 + \sum_{i=1}^{3}\pi_i^2\right)^2 + \sum_{i=1}^{3}\pi_i(\Box\pi_i) + \sigma(\Box\sigma) \tag{26}$$

instead of eq. (21) that is obtained from eq. (26) with the help of equations of motion. In the chiral limit the only nonzero contribution to the matrix element (25) is from the last term in eq. (26), and it immediately gives the correct result (eq. (25)).

III. CONCLUSION

In the present paper we have proved eq. (1) in a new way using effective linear as well as nonlinear Lagrangian models. We have seen that while in the case of linear σ-model the validity of eq. (1) is not so obvious at first sight, in the case of nonlinear phenomenological Lagrangians the validity of eq. (1) is evident just from eqs. (6) and (12) (or (13)). We have also derived a set of new matrix elements (see, e. g., eqs. (17) and (18)) in which the states are not only with physical masses but also from the whole pseudoscalar meson nonet, including the η' particle and correct $\eta\eta'$ mixing.

All these relations (eqs. (1) and (18)) are examples of strong, unsuppressed transitions between quark and gluon degrees of freedom in the 0^+ channel, and such relations have to be satisfied [14, 20–21] by any realistic model of coupling between the scalar gluonium and pseudoscalar mesons.

In this connection we believe that especially eqs. (18) could be helpful in understanding the recently discovered scalar meson G (1590) [22]. Having dominant decays into $\eta\eta$ and $\eta\eta'$ channels [22] this meson has been interpreted [23] as a pure scalar gluonium. However, this interpretation a priori assumes strong suppression of coupling between gluons and quarks (and hence, a necessary suppression of decays into $\pi\pi$ and $K\bar{K}$) in the 0^+ channel; the assumption is model-dependent (based on large N_C dynamics) and has been shown here not to be true. In fact, there are models of a scalar glueball [20–21], and these models (respecting eq. (1)) lead to a large width for a heavy scalar gluonium decaying into two pions [20–21]. On the other hand, supposing consistently with eq. (1) large and unsuppressed transitions between quarks and gluons in the 0^+ channel we have shown [24] that the G (1590) meson can be interpreted as the unitary singlet scalar quarkonium or its radial excitation, too (see also [25]). Even with these interpretations it could not be so surprising that the decay G (1590) $\to \pi\pi$ is strongly suppressed because there exist examples of such suppressions in the hadronic world, e. g., while $\Gamma_\varrho \approx \Gamma(\varrho \to \pi\pi) = 154$ MeV [5] its radial excitation ϱ' (1220) is strongly suppressed to decay into $\pi\pi$, i. e. $\Gamma(\varrho' \to \pi\pi) \approx 3$ MeV but the total width $\Gamma_{\varrho'} = (200 \div 300)$ MeV [26]. Moreover, if some wide structure in the G (1590) region in the $\pi\pi$ system exists [27], then the interpretation of the G (1590) meson could be even more complicated.

We would like to thank Drs. A. T. Filippov, A. B. Govorkov, A. K. Likhoded for discussions and Prof. V. A. Meshcheryakov for interest and support in this work.

REFERENCES

[1] Novikov, V. A. et al.: Nucl. Phys. B *191* (1981), 301.
[2] Zweig, G.: *CERN preprints* TH-401, TH–412 (1964);
Okubo, S.: Phys. Lett. *5* (1963), 165;
Iizuka, J.: Prog. Theor. Phys. Suppl. *37—38* (1966), 21.
[3] t'Hooft, G.: Nucl. Phys. B *72* (1974), 461.
[4] Robson, D.: Nucl. Phys. B *130* (1977), 328.
[5] Particle Data Group: Rev. Mod. Phys. *56* (1984), S1.
[6] Novikov, V. et al.: Nucl. Phys. B *165* (1980), 55.
[7] Voloshin, M., Zakharov, V.: Phys. Rev. Lett. *45* (1980), 688;
Novikov, V., Shifman, M.: Z. Phys. C *8* (1981), 43.
[8] Gasiorowicz, S., Geffen, D. A.: Rev. Mod. Phys. *41* (1969), 531.
[9] Witten, E.: Annals of Phys. *128* (1980), 363;
di Vecchia, P. et al.: Nucl. Phys. B *181* (1981), 318;
Rosenzweig, C., Schechter, J., Trahern, C. G.: Phys. Rev. D *21* (1980), 3388.

[10] Collins, J. C., Duncan, A., Joglekar, S. D.: Phys. Rev. D *16*(1977), 438.
[11] Callan C. G., Coleman, S., Jackiw, R.: Annals of Phys. *59* (1970), 42.
[12] Lánik, J.: JINR E2-84-437, Dubna, 1984 and Talk given at the *Internat. School of Elementary Particle Physics*, Kupari-Dubrovnik, Yugoslavia, 1984;
Lánik, J.: Acta Phys. Slov. to be published.
[13] Ellis, J.: Nucl. Phys. B *22*(1970), 478.
[14] Lánik, J.: Phys. Lett. *144 B*(1984), 439.
[15] Shifman, M. A., Vainshtein, A. I., Zakharov, V. I.: Nucl. Phys. B *147*(1979), 385, 448.
[16] Bell, J. S., Bertlmann, R. A.: Nucl. Phys. B *177*(1981), 218.
[17] Kawarabayashi, K., Ohta, N.: Nucl. Phys. B *175*(1980), 477.
[18] Filippov, A. T.; Uspekhi Fiz. Nauk *137*(1982), 201;
Gasser, J., Leutwyler, H.: Nucl. Phys. B *250*(1985), 465.
[19] JADE Collab., Bartel, W. et al.: Phys. Lett. *158*(1985), 511.
[20] Gomm, h., Schechter, J.: Phys. Lett. *158 B*(1985), 448.
[21] Ellis, J., Lánik, J.: Phys. Lett. *150 B*(1985), 289;
Lánik, J.: Paper contributed to the *Internat. Europhys. Conf. on High Energy Physics*, Bari 1985 and Dubna report E2-85-260, 1985.
[22] Binon, F. et al.: Nuovo Cim. *78 A*(1983), 313;
ibid *80 A*(1984), 365.
[23] Gershtein, S. S., Likhoded, A. K., Prokoshkin, Yu. D.: Z. Phys. C Particle and fields *24* (1984), 305.
[24] Lánik, J.: Pisma v ZhETF *42*(1985), 122.
[25] Mennessier, G., Narison, S., Paver, N.: Phys. Lett. *158 B*(1985), 153;
Govorkov, A. B.: JINR, P2-85-696, Dubna 1985.
[26] Gerasimov, S. B., Govorkov, A. B.: Z. Phys. C — Particle and Fields *13*(1982), 43;
Dubnička, S., Martinovič, Ľ.: in preparation.
[27] Krupa, D., Meshcheryakov, V. A., Surovcev, Yu. S.: JINR, P2-84-857, Dubna 1984.

Received November 5th, 1985