ITEP – 9

INSTITUTE OF THEORETICAL
AND EXPERIMENTAL PHYSICS

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TOWARDS A THEORY OF WEAK HADRONIC DECAYS OF CHARMED PARTICLES

MOSCOW 1986

УДК 539.172 М-16

A b s t r a c t

Weak decays of charmed mesons are considered. We propose a new quantitative framework for theoretical analysis of nonleptonic two-body decays based on the QCD sum rules. This is the first of a series of papers devoted to the subject. We discuss theoretical foundations of the approach ensuring model-independent predictions for the partial decay widths.

Submitted to Yadernaya Fizika.

1. Introduction

Starting from 1976 when charmed particles were discovered[1] and till present days the weak exclusive decays of charmed mesons attract interest of both theorists and experimentalists. Experimentally, large number of partial decay widths has been measured, and the efforts in this direction continue. Thus, a recent experiment[2] of MARK III group has yielded a set of data essentially differing from previous results[3]. Comparing [2] and [3] and results of other groups (see, e.g., the review[4]) we see that the situation is not quite settled, the numbers breath. At the same time one can hardly expect drastic changes in the general picture comprising ~ 50 decay channels.

With such an array of apparently disconnected data the need in theoretical interpretation becomes extremely acute. Are we able to understand the results theoretically? What is even more important, can we learn something new about weak or strong interactions? Ever since 1976 the weak exclusive decays of D- and F-mesons have been a continuous challengeto the theory.

We will try to demonstrate that the answer to the first question is positive. In this paper we develop an approach allowing one to analyse the majority of D- and F-meson decays in two pseudoscalars and one pseudoscalar plus one vector. (The set of decay modes to be considered is given in the Table, altogether about 50 modes). Our approach is based on QCD sum rules[5]. The predictions for the partial widths will be practically model-independent (the estimated uncertainty in decay rates reduces to a factor ~ 1.5).

As for the second question, unfortunately, we must confess that the analysis adds rather little to what is already known in

QCD. Starting from the already developed picture of the QCD vacuum we simply explain peculiarities of the exclusive decays, reproducing experimental numbers with reasonable accuracy. Perhaps, the most important lesson is confirmation of the essential role of soft gluons, the observation made previously (see, e.g. [17, 19]).

We start from a brief historical survey which will introduce us into a theoretical "kitchen" and will serve as a reference point in our constructions.

The first works (e,g., ref. [6]) published even prior to discovery of open charm exploited, in essence, a naive spectator model; at those days the charmed-quark mass m_c was believed to be large enough. It became clear soon, however, that the corresponding predictions were in sharp disagreement with the data (see, e.g. Sect.3.5 of the review paper [7]), and pre-asymptotic effects of different nature must play an essential role.

There exists a variety of models intended for the description of the weak non-leptonic decays of D- and F-mesons. They can be grouped in a few classes by the preference given to this or that preasymptotic effect. Let us catalogue here some popular schemes: (a) Spectator mechanism [6,8]; (b) QCD corrections to the coefficients in H_W beyond the leading log approximation [9]; (c) Final state interactions [10]; (d) Non-spectator diagram dominance (annihilation mechanism) [11-15]; (e) Emission of perturbative gluons [16-19]; (f) Attempts to account for non-perturbative gluons [17,20,21,40]; (g) Phenomenological modifications of the spectator model [22,23]; (h) Interference effects [24,41,42]; (i) Low-energy lagrangians for mesonoc fields (e.g., [25,43]); (j) SU(3) symmetry predictions, in particular, sextet dominance [26-29]; (k) Fitting the mixing angles in the Kobayashi-Maskawa matrix [30] (for experimental values see, however, refs. [46-48]); (l) Penguin diagram dominance [31,32,45]; (m) Constituent versus current quarks [33]; (n) $1/N_c$ expansion [23, 24].

The naive spectator model - historically the first - emerged as a natural successor to the model used in the strange particle decays. It assumes that D or F decay proceeds as if the c-quark is free and the transition is not affected by the light spectator quark. The width of the two-particle decay is determined by the graph of Fig.1a which is treated within the factorization hypothesis. To illustrate the approach, consider $D^+ \to \bar{K}^0 \pi^+$ channel. The amplitude is given by $\langle D^+ | H_W | \bar{K}^0 \pi^+ \rangle$ where H_W is the effective $\Delta C = 1$ hamiltonian [35] including the hard gluon exchanges; we assume standard GIM scheme [44] for s-d mixing and omit b, t effects:

$$H_W = \frac{1}{\sqrt{2}} G_F \left(C_1 \bar{s}' \gamma_\mu (1+\gamma_5) c \, \bar{u} \gamma_\mu (1+\gamma_5) d' + C_2 \bar{s}' \gamma_\mu (1+\gamma_5) d' \, \bar{u} \gamma_\mu (1+\gamma_5) c \right)$$

$$d' = d \cos\theta + s \sin\theta , \quad s' = s \cos\theta - d \sin\theta \tag{1}$$

$$C_1 = \frac{1}{2}(C_+ + C_-) , \quad C_2 = \frac{1}{2}(C_+ - C_-)$$

$$C_- = \left(\frac{\alpha_s(m_c)}{\alpha_s(m_W)} \right)^{4/b} \simeq 1.57 , \quad C_+ = \frac{1}{\sqrt{C_-}} \simeq 0.81 , \; C_1 = 1.19 , \; C_2 = -0.38 \tag{2}$$

Here m_c = 1.35 GeV, m_W = 80 GeV are the masses of the c-quark and W boson, θ = 13° is the Cabibbo angle, G_F is the Fermi coupling constant, b = 25/3 is the coefficient in the Gell-Mann-Low function,

$$\alpha_s(Q^2) = \frac{4\pi}{b \ln Q^2/\Lambda^2} \tag{3}$$

For the QCD scale parameter Λ in eq.(3) we accept Λ = 100 MeV, in accordance with the modern data [36]. (Notice that in previous analyses of weak decays a much larger value, $\Lambda \sim 500$ MeV, was popular); For the $D^+ \to \bar{K}^0 \pi^+$ amplitude we get

$$M(D^+\to \bar{K}^0\pi^+) = \frac{1}{\sqrt{2}} G_F \cos^2\theta \{ c_1 \langle \bar{K}^0 | \bar{s}^i \gamma_\alpha (1+\gamma_5) c^i | D^+ \rangle \langle \pi^+ | \bar{u}^j \gamma_\alpha (1+\gamma_5) d^j | 0 \rangle +$$
$$+ c_1 \langle \pi^+ | \bar{u}^i \gamma_\alpha (1+\gamma_5) c^i | D^+ \rangle \langle \bar{K}^0 | \bar{s}^i \gamma_\alpha (1+\gamma_5) d^j | 0 \rangle +$$
$$+ c_2 \langle \bar{K}^0 | \bar{s}^i \gamma_\alpha (1+\gamma_5) c^j | D^+ \rangle \langle \pi^+ | \bar{u}^j \gamma_\alpha (1+\gamma_5) d^i | 0 \rangle +$$
$$+ c_2 \langle \pi^+ | \bar{u}^i \gamma_\alpha (1+\gamma_5) c^i | D^+ \rangle \langle \bar{K}^0 | \bar{s}^j \gamma_\alpha (1+\gamma_5) d^j | 0 \rangle = \tag{4}$$
$$= \frac{1}{\sqrt{2}} G_F \cos^2\theta \{ (c_1 + c_2/3) f_+^K f_\pi + (c_2 + c_1/3) f_+^\pi f_K \} (-i m_D^2)$$

where m_D is the D^+ mass,

$$\langle \pi^+(p) | \bar{u}^i \gamma_\alpha \gamma_5 d^j | 0 \rangle = -\frac{1}{3} i \delta^{ij} f_\pi p^\alpha , \tag{5}$$

$$\langle \bar{K}^0(p^K) | \bar{s}^i \gamma_\alpha c^j | D^+(p) \rangle = \frac{1}{3} \delta^{ij} [f_+^K (p+p^K)_\alpha + f_-^K (p-p^K)_\alpha] . \tag{6}$$

Neglecting the SU(3) symmetry breaking and light-meson masses (in this limit $f_\pi = f_K = 133$ MeV) we arrive at

$$M(D^+ \to \bar{K}^0 \pi^+) = -\frac{i m_D^2}{\sqrt{2}} G_F \cos^2\theta f_+^K f_\pi \frac{4}{3} (c_1 + c_2) \tag{7}$$

The constant f_+^K has been determined by virtue of the QCD sum rules in ref.[37], $f_+^K = 0.5 \pm 0.1$. Then

$$Br(D^+ \to \bar{K}^0 \pi^+) \sim 6\% , \tag{8}$$

a factor 2.5 higher than the experimental number. Even a larger discrepancy between the spectator model and experiment emerges in the ratio

$$\alpha_{theor} = \frac{\Gamma(D^0 \to \bar{K}^0 \pi^0)}{\Gamma(D^0 \to K^- \pi^+)} = \frac{1}{2} \cdot \frac{1}{9} \frac{(2c_+ - c_-)^2}{(\frac{2}{3} c_+ + c_-)^2} \sim 10^{-3} \div 10^{-2} \quad (9)$$

(while $\alpha_{exp} \simeq 0.35 \pm 0.1$). To eliminate the discrepancy some authors invoke final state interactions (e.g., [19]) which, on the contrary, overestimates the ratio, $\alpha_{final\ state\ int} \sim 2$. The recent discovery of the $D^0 \to \bar{K}^0 \phi$ decay with 1% branching ratio [38] is also in a sharp contradiction with the spectator model yielding for this branching ratio $\sim 10^{-2}$%. Thus, the spectator model and its simple modifications cannot describe two-particle decays of D and F mesons even in a rough approximation.

Still, we dwell here on the spectator model keeping in mind that below we will use it as a reference point. The spectator model (to be referred below as the standard model) gives a convenient and relatively simple and transparent language, since it allows one to discuss other models and mechanisms in terms of deviations from the standard model. Just these deviations will be the main subject of analysis in the present work.

Notice that the factorization property in the matrix element like $\langle 2\pi | H_w | K \rangle$ is intensively exploited in the theory of non-leptonic decays of strange particles [39] where it seems to work with a satisfactory accuracy. In description of the strange particle decays factorization is, at least, compatible with the soft pion technique [39]; this fact might serve as a justification. Analogous arguments based on the small energy release and PCAC show that one can neglect the final state interactions in the kaon and hyperon non-leptonic decays. For charmed mesons the energy release is relatively large ($\sim$1 GeV) and these approximations are nc more reliable.

Critical analysis of all later models aimed at improving the standard model is beyond the scope of this paper. The general conclusion can be summarized in a few words: soft gluons (and/or quark pairs) should play an essential role in the two-particle amplitudes. The inclusive D and F decays are not exhausted by simplest quark graphs. Especially instructive in this respect is ref.[19] which stimulated much works in the field. Unfortunately, in the quantitative aspect the approaches developed so far suffer from a common drawback: the absence of the direct relation to the fundamental QCD. Some of them are motivated by QCD but still use additional assumptions. Therefore, one cannot say beforhand whether or not approximations made are reliable, whether or not the accepted values of parameters are reasonable, one cannot estimate theoretically (with no comparison with experiment) accuracy of the predictions.

Our aim is to develop a new approach tending to avoid model dependence or, at least, minimize it drastically. As has been already mentioned, the approach is based on the QCD sum rules[5] (for an extensive review see[50]); conceptually it is close in spirit to the method used in ref.[55] for the analysis of hyperon decays. It will allow us to consider within a unified scheme all transitions

$$D \to PP \tag{10}$$

$$D \to PV \tag{11}$$

$$F \to PP \tag{12}$$

$$F \to PV \tag{13}$$

where $P = \pi, K, \eta$ stand for the Goldstone meson and V is a generic notation for vector mesons, ρ, ϕ, ω, K^*.

The paper is organized as follows. In Sec.2 we formulate the problem in terms suitable for the sum rule method. In Sec.3 a general form of the sum rules for weak two-partical amplitudes is established. Sec.4 demonstrates that the sum rules can be written down directly for the difference $M-M_f$, where M is the total amplitude while M_f refers to the factorizable part. Basic results are summarized in Sec.5.

2. Outlining the Problem

The following four-point amplitude will play the basic role in our analysis:

$$\Pi_{\mu\nu}(Q_1, Q_2; q) = \int d^4x\, d^4y\, d^4z\, e^{iQ_2x + iqy} \langle T\{j_{A_\mu}(y), j_{B_\nu}(x), j_D(0), H_W(z)\}\rangle_0 \tag{14}$$

Here j_D, j_A, j_B are the quark currents producing the corresponding mesons from the vacuum, for instance

$$j_{D^0} = i\bar{c}\gamma_5 u, \quad j_{D^+} = i\bar{c}\gamma_5 d \tag{15}$$

$$j_{F^+} = i\bar{c}\gamma_5 s \tag{16}$$

H_W is the weak $\Delta C=1$ hamiltonian (see eq.(1)). The four-point function (14) is depicted in Fig.2 where the shaded blob denotes the weak decay amplitude of the charmed meson into two light mesons - A and B.

To avoid direct instantons [52] we use for light pseudoscalars axial-vector, not pseudoscalar currents,

$$j_{\pi^+} = \bar{d}\gamma_\lambda\gamma_5 u,\ j_{\pi^0} = \frac{1}{\sqrt{2}}(\bar{u}\gamma_\lambda\gamma_5 u - \bar{d}\gamma_\lambda\gamma_5 d),\ j_{\pi^-} = \bar{u}\gamma_\lambda\gamma_5 d \tag{17}$$

for pions, and

$$j_{K^+} = \bar{s}\gamma_\lambda\gamma_5 u,\ j_{\bar{K}^0} = \bar{s}\gamma_\lambda\gamma_5 d,\ j_{K^-} = \bar{u}\gamma_\lambda\gamma_5 s \tag{18}$$

for kaons. This has another advantage - all amplitudes (10)-(13) can actually be treated simultaneously. A small octet-singlet mixing for pseudoscalars (i.e. $\eta_1 - \eta_8$) is neglected, so that the currents producing η, η' are

$$j_\eta = \frac{\bar{u}\gamma_\lambda\gamma_5 u + \bar{d}\gamma_\lambda\gamma_5 d - 2\bar{s}\gamma_\lambda\gamma_5 s}{\sqrt{6}},\ j_{\eta'} = \frac{\bar{u}\gamma_\lambda\gamma_5 u + \bar{d}\gamma_\lambda\gamma_5 d + \bar{s}\gamma_\lambda\gamma_5 s}{\sqrt{3}} \tag{19}$$

For vector mesons, ρ, K^*, ω, ϕ we use

$$j_{\rho^-} = \bar{u}\gamma_\lambda d,\ j_{\rho^0} = \frac{1}{\sqrt{2}}(\bar{u}\gamma_\lambda u - \bar{d}\gamma_\lambda d),\ j_{\rho^+} = \bar{d}\gamma_\lambda u \tag{20}$$

$$j_{K^{*+}}=\bar{s}\gamma_\mu u,\; j_{\bar{K}^{*0}}=\bar{s}\gamma_\mu d,\; j_{K^{-*}}=\bar{u}\gamma_\mu s \tag{21}$$

$$j_\omega=\frac{1}{\sqrt{2}}(\bar{u}\gamma_\mu u+\bar{d}\gamma_\mu d) \tag{22}$$

$$j_\phi=\bar{s}\gamma_\mu s \tag{23}$$

The chiral limit, $m_u=m_d=m_s=0$ is assumed throughout the work. This assumption simplifies computations. All Goldstone mesons, π , K , η , have vanishing masses and their residues are determined by a common constant, f_π = 133 MeV. In the chiral limit all vector mesons also have a common mass, coinciding with that of the ρ , m_ρ = 0.78 GeV. Their residues can be expressed in terms of f_ρ by virtue of the SU(3)$_{flavour}$ symmetry. Here

$$\langle\rho^-(p)|\bar{u}\gamma_\mu d|0\rangle=f_\rho m_\rho e_\mu(p); \tag{24}$$

numerically $f_\rho\simeq 200$ MeV . (Experimentally f_ρ is measurable in $\rho^0\to e^+e^-$). Furthermore, e_μ is the ρ-meson polarization vector. In the case of η' meson we use $f_{\eta'}\simeq 100$ MeV [53], $m_{\eta'}$ = 0.96 GeV (in this channel there is no massless particle in the chiral limit).

As has been noted above, a characteristic feature of our approach is the possibility to formulate predictions for the amplitudes directly for the deviations from the factorized values. A systematic discussion of this point will be given later, and herewe only sketch the basic idea.

There are three "skeleton" diagrams which determine the four-point function (14). They are depicted in Figs.3a,b,c. (Actual graphs are obtained from the skeleton diagrams but cutting off some solid line and/or adding gluon exchange (see, e.g. Fig.3d).

The graphs in Fig.3b,c correspond to the spectator mechanism (factorization is not assumed, however) while the diagram 3a corresponds to the so called annihilation mechanism.

Consider, for instance, the diagram 3d which, within our approach, is relevant to the sum rules for the $D^0 \to \bar{K}^0 \phi$ decay. In the quark language one can interprete the graph as follows: in the initial D^0 meson $\bar{u}$ quark emits a soft (non-perturbative) gluon; as a result of weak interaction there emerge two light quarks, s and d, one of them annihilates the soft gluon emitted by $\bar{u}$; the final quarks combine to produce K and ϕ. We treat the diagram by virtue of a "refined duality" of the QCD sum rules, and in this way obtain a calculational scheme allowing us to reliably estimate (under theoretical control) the effects of soft gluon (or quark) emission. The scheme is very similar to the original sum rules[5]. The only non-technical additional difficulty is a contamination of the phenomenological part of the sum rules by some "parasitic" contributions which enter with no exponential suppression (see below).

Relative weights of diagrams 3a,b,c are different in each particular decay mode, and combination of these factors results in a rich and peculiar hierarchy of amplitudes.

Let us discuss now to which extent the approximations accepted above are adequate.

Chiral corrections due to $m_s \neq 0$ and $\langle\bar{s}s\rangle_0 - \langle\bar{u}u\rangle_0 \neq 0$ are of two types. In the theoretical part they are of order $m_s/m_c \sim 10^{-1}$. Besides that, in the phenomenological part one must substitute

$$\dagger \; \frac{1}{f_\pi} \to \frac{1}{f_K} \exp(m_K^2/M'^2), \tag{25}$$

$$\frac{1}{f_\rho} e^{m_\rho^2/M^2} \to \frac{1}{f_{K^*}} \exp(m_{K^*}^2/M'^2). \tag{26}$$

As we shall see below, the characteristic value of M'^2 is of order 1 GeV^2 and both effects are expected to be $\lesssim 10\%$ at least in the process with one s quark in the final state. We will neglect them. With two s-quarks in the final state chiral corrections may be a factor of two larger. $SU(3)_{fl}$ violating effects may be most noticeable in the phase space, and here we, of course, will take them into account.

3. Description of the Method

Following the standard procedure[5] we write down a double dispersion relation for each invariant amplitude in the four-point function (14):

$$f(Q_1^2, Q_2^2, q^2) = \iint \frac{\rho(s,s',q^2)\,ds'\,ds}{(s-Q_1^2)(s'-Q_2^2)} + \text{(subtraction terms)} \tag{27}$$

Here Q_1, Q_2 and q are the momenta of the virtual D meson and quark currents producing D and A mesons, respectively.

Furthermore, the phenomenological part (r.h.s.) of the sum rules (27) is saturated by resonances, (while the theoretical part is calculated with the aid of OPE in the domain $Q_1^2, Q_2^2, q^2 \sim -1\ \mathrm{GeV}^2$, when, on one hand, QCD is still applicable, and on the other hand, effects due to non-trivial vacuum structure are quite noticeable.

Technically, the procedure is more labour-consuming than in popular model formulated directly in physical (Minkowski) space. Indeed, we work in the euclidean domain of momenta, and have to put additional efforts to extrapolate the results to the physical domain. The procedure turns out to be less transparent. As a reward, we get the possibility of reliable calculations, not just rough estimate, of non-perturbative effects.

As we shall see below, the characteristic values of momenta, both external and virtual, are about -1 GeV^2 . This means that we can indeed use the effective hamiltonian (1) with only hard gluon exchanges included C_+, C_- . The issue of relatively soft gluons in H_W does not arise, and we can treat C_+, C_- as constants indicated in eq.(2).

The next step is the choice of an appropriate kinematic structure in the four-point function (14). The general decomposition of $\Pi_{\mu\nu}$ is as follows

$$\Pi_{\mu\nu}(Q_1, Q_2; q) = f_1 g_{\mu\nu} + f_2 q^\mu Q_2^\nu + f_3 q^\nu Q_2^\mu + f_4 q^\mu q^\nu + f_5 Q_2^\mu Q_2^\nu \tag{28}$$

(Here we have taken into account that $Q_1 = Q_2 + q$, so that there are only two independent momenta). In the case of $D \to PP$ and $F \to PP$ the relevant structure is, evidently, $q^\mu Q_2^\nu$ since the matrix elements $\langle 0 | j_{A,B} | P \rangle$ are proportional to the momenta of the pseudoscalar mesons at hand. For the decays $D \to PV$, $F \to PV$ the relevant structure can be found from two requirements; first, it should contain the pseudoscalar meson momentum q^μ ; second, it should be in correspondence with the form of the amplitude

$$\mathcal{M}(D \to PV) = Ceq \tag{29}$$

where e is the polarization vector of V, C is some constant. As a consequence, the structure of interest is

$$T_{\gamma\nu} = q^{\gamma}\left(q^{\nu} - \frac{qQ_2}{Q_2^2} Q_2^{\nu}\right) g(Q_1^2, Q_2^2; q^2) \tag{30}$$

where g is the invariant amplitude.

On the other hand, $\Pi_{\gamma\nu}$ defined in eq.(14) must be transversal

$$Q_2^{\nu} \Pi_{\gamma\nu} = 0 \tag{31}$$

up to commutator terms which are irrelevant, however, to vector-meson contribution (the commutator terms correspond to those pieces in the r.h.s., which have no vector-meson pole. This contribution in the sum rules is expected to be negligible; see below). Therefore, it is sufficient to consider just the same structure, $q^{\gamma} Q_2^{\nu}$, as in the D → PP case. Knowing the coefficient in front of $q^{\gamma} Q_2^{\nu}$ we readily reconstruct g figuring in eq.(30) by virtue of the condition (31). All further analysis will refer just to this structure.

Discussion of calculational details in the theoretical part is suspended till the next paper. Here we dwell on two general remarks.

In the operator expansion we will keep only operators with dimension $d \leq 6$, resulting in the following v.e.v.'s: $\langle G^2 \rangle_0$ $\langle \bar{\psi}\psi \rangle_0$, $\langle \bar{\psi} G_{ij} \sigma_{ij} \psi \rangle_0$, $\alpha_s \langle \bar{\psi}\psi \rangle_0^2$, and the unit operator I. Dim 7 and higher operators in our conditions are inessential - this can be checked by an explicit estimate. All coefficients in the OPE will be calculated in the leading (non-vanishing) order in α_s.

The version of the sum rules we will use is close in spirit to the variant exploited in refs. [51,54,55] for determination of baryon magnetic moments, charge radii and other analogous characteristics. Indeed, in the both cases one must switch on an external (with respect to QCD) interaction - electromagnetic in ref. [54] and H_W in the case at hand. Technically, the treatment of weak decays of charmed mesons is more complicated than, say, the problem of the pion charge radius, since the number of the particles involved in the process is larger. Fortunately, one aspect is, on the contrary, simpler. Unlike ref. [51,54,55] there is no need for us to introduce new (induced) vacuum condensates (for instance, in ref. [54] a vacuum magnetic susceptibility in the external electromagnetic field has been introduced).

Let us elucidate the assertion in more detail. Fig.4a displays the amplitude for the theoretical part of the sum rules in the problem of the pion charge radius. The momentum flowing in the electromagnetic current is assumed to be small. Formally, an external field $F_{\mu\nu}$ is introduced and the correlator $\langle T\{j_\pi(x), j_\pi(0)\}\rangle_0$ is considered in the linear in $F_{\mu\nu}$ approximation. In this kinematics the quark lines marked by shadded blobls can be almost on the would-be mass shell. We are unable to calculate the corresponding effect and parametrize it by a new v.e.v., $\langle\bar\psi\sigma_{\mu\nu}\psi\rangle_{F_{\mu\nu}}$, cutting off all soft lines (Fig.4b).

It seems at first sight that the situation with weak decays is quite analogous. The weak hamiltonian brings no momentum (Figs. 2,3). This means that all lines attached to H_W can be "soft" in the sense that $p\sim 0$. Then light quarks are indeed close to the mass shell; however, the off -shellness of the c-quark is large, $m_c^2 \gg \Lambda^2$. Therefore, with respect to the c-quark the perturbative decsription is applicable and hence there is no disconnected block, isolated from the hard part of the diagram (as in Fig.4b). The corresponding non-perturbative contribution is depicted, for instance, in Fig.4c - it reduces to $\langle(\bar\psi\psi)^3\rangle_0$, no new v.e.v. is necessary.

In this respect the analysis of weak decays of, say, kaons would be different, since for kaons (if they are treated within the sum rule method) one could not avoid introducing new induced v.e.v.'s reflecting soft interactions with an auxiliary external field, h, $\mathcal{L}_{int} = h H_W$.

Let us turn now to the phenomenological part of the sum rules. It can be represented as a sum of diagrams depicted in Figs.5a-d. Various pieces here describe different options. Fig.5a gives a three-resonance piece containing the decay amplitude of interest (we denote it by T)

$$F_1(Q_1^2, Q_2^2; q^2) = \frac{T f_A f_B f_D\, m_D^2/m_c}{q^2(Q_1^2-m_D^2)(Q_2^2-m_B^2)} + \sum_k\sum_n\sum_j \frac{T_{knj}\, f_i f_n f_j}{(Q_1^2-m_k^2)(Q_2^2-m_n^2)(q^2-m_j^2)} \qquad (32)$$

Here $\langle 0|\bar c\gamma_5 u|D^0\rangle = i f_D m_D^2/m_c$, f_D = 170 MeV [49], and f_A, f_B stand for the residues of A and B into the corresponding currents, m_B is the B particle mass (m_A = 0 in the chiral limit). The second term in eq.(32) corresponds to higher-resonance contribution. More exactly, its interpretation is: the current j_D produces the k-th excitation which decays to j-th and n-th states annihilated by the currents j_A, j_B .

Apart from the diagram 5a in saturating the r.h.s. one must take into account the graphs depicted in Fig. 5b,c,d. They yield one-pole terms describing non-diagonal transitions in the sum rules for correlation functions in the external field [53-55], see also Fig.6 where we have displayed for comparison analogous graphs for the proton-proton transition in the axial field). Figs.5 b-d represent a contamination, a "parasitic" background with respect to the amplitude of interest (Fig.5a). Our method will be reliable only if the background is not large, say, does not exceed 20-30%. Then, even a rough estimate of the background contribution, valid up to a factor of 2, will result in a tolerable error ($\sim$ 10%) in the estimates of the decay amplitudes.

For an approximate estimate of the contamination due to the diagrams 5 b-d we will invoke a simplest pole model (Figs. 7a,b). Then we get for the diagram 5b

$$F_2(Q_1^2,Q_2^2,q^2)=\sum_j \frac{T_{Dj}\,T_{jAB}\,f_A f_B f_D\, m_D^2/m_c}{(Q_1^2-m_D^2)(Q_2^2-m_B^2)(Q_1^2-m_j^2)q^2}+ \tag{33}$$
$$+\sum_{i,j,k,s}\frac{T_{ij}\,T_{jks}\,f_i f_k f_s}{(Q_1^2-m_i^2)(Q_2^2-m_k^2)(Q_1^2-m_j^2)(q^2-m_s^2)}$$

Here j is an intermediate state contaning no heavy quark and decaying to A and B by virtue of the strong interaction, $T_{ij}=\langle i|H_W|j\rangle$ $T_{Dj}=\langle D|H_W|j\rangle$; T_{jks} is the strong decay amplitude $j\to ks$ m_j, m_k, m_s denote the masses of the corresponding states. The diagram 5c reduces to

$$F_3(Q_1^2,Q_2^2,q^2)=\sum_j \frac{T_{DjA}\,T_{jB}\,f_A f_B f_D\, m_D^2/m_c}{(Q_1^2-m_D^2)(Q_2^2-m_B^2)(Q_2^2-m_j^2)q^2}+ \tag{34}$$
$$+\sum_{i,j,k,s}\frac{T_{isj}\,T_{jk}\,f_i f_k f_s}{(Q_1^2-m_i^2)(Q_2^2-m_j^2)(Q_2^2-m_k^2)(q^2-m_s^2)}$$

Here $T_{jk}=\langle j|H_W|k\rangle$, $T_{jB}=\langle j|H_W|B\rangle$; T_{isj} is the amplitude of the strong transition of the state i (produced by j_D) to a state j (containing c-quark) and a light particle s.

In eqs.(33),(34) the pieces corresponding to the transition $D \to AB$ and to those with higher resonances are written out separately.

The basic sum rule takes the form

$$f_T(Q_1^2, Q_2^2, q^2) = F_1(Q_1^2, Q_2^2, q^2) + F_2(Q_1^2, Q_2^2, q^2) + F_3(Q_1^2, Q_2^2, q^2) + F_4(Q_1^2, Q_2^2, q^2) \quad (35)$$

where f_T is the theoretical part, F_i (i = 1-4) are the functions which come from Figs.5 a-d, respectively. Fig.5d does not depend on Q_1^2 and Q_2^2 simultaneously.

Let us proceed now in the euclidean domain, $Q_1^2 \to -Q_1^2, Q_2^2 \to -Q_2^2$, $q^2 \to -q^2 \equiv Q^2$. To suppress higher resonances in the r.h.s. of eq.(35) we apply the double Borel transformation, as in ref.[56]

$$\hat{B} f(Q_1^2, Q_2^2, Q^2) = \lim_{\substack{n', n \to \infty \\ Q_1^2, Q_2^2 \to \infty \\ Q_1^2/n = M^2, Q_2^2/n' = M'^2}} \frac{M^2 M'^2}{(n-1)!(n'-1)!} (Q_1^2)^n \left(\frac{d}{dQ_1^2}\right)^n \times \quad (36)$$

$$\times (Q_2^2)^{n'} \left(\frac{d}{dQ_2^2}\right)^{n'} f(Q_1^2, Q_2^2, Q^2) = f_B(M^2, M'^2, Q^2) .$$

It is instructive to notice that the spectral density in eq.(28) is actually a sum of spectral densities emerging from graphs of Fig.5 which results in a variety of thresholds and different continuum structures in the sum rules.

The double Borel transformation kills the one-pole contribution of Fig.5d. For the function F_1 we obtain

$$F_1(M^2, M'^2, Q^2) = \frac{T f_A f_B f_D m_D^2 / m_c}{} e^{-m_D^2/M^2 - m_B^2/M'^2} +$$

$$+ \sum_k \sum_n \sum_j \frac{T_{knj} f_i f_n f_j}{(Q^2 + m_j^2)} e^{-m_k^2/M^2} e^{-m_n^2/M'^2} \quad (37)$$

The Borel transform of F_2 takes the form

$$F_2(M^2, M'^2, Q^2) = \sum_j \frac{T_{Dj} T_{jAB} f_A f_B f_D \, m_D^2/m_c}{Q^2} \frac{e^{-m_B^2/M'^2}(e^{-m_D^2/M^2} - e^{-m_j^2/M^2})}{m_D^2 - m_j^2} + \tag{38}$$
$$+ \sum_{i,j,k,S} \frac{T_{ij} T_{jkS} f_i f_k f_S}{(Q^2 + m_S^2)} \frac{e^{-m_k^2/M'^2}(e^{-m_i^2/M^2} - e^{-m_j^2/M^2})}{m_i^2 - m_j^2}$$

And, finally, the Borel transform of F_3 is

$$F_3(M^2, M'^2, Q^2) = \sum_j \frac{T_{DjA} T_{jB} f_A f_B f_D \, m_D^2/m_c}{Q^2} \frac{e^{-m_D^2/M^2}(e^{-m_B^2/M'^2} - e^{-m_j^2/M'^2})}{m_B^2 - m_j^2} + \tag{39}$$
$$+ \sum_{i,j,k,S} \frac{T_{iSj} T_{jk} f_i f_k f_S \, e^{-m_i^2/M^2}}{(Q^2 + m_S^2)} \frac{(e^{-m_j^2/M'^2} - e^{-m_k^2/M'^2})}{m_j^2 - m_k^2}$$

One can easily see that each of the terms in eqs.(38),(39) is a sum of two pieces. The first piece is proportional to $exp(-m_D^2/M^2)\cdot exp(-m_B^2/M'^2)$ (it emerges after contraction of the pole due to the intermediate state, in Fig.5c). Quite evidently, this piece merely renormalizes the genuine weak decay amplitude. The second piece, just the most dangerous, contains $exp(-m_i^2/M^2)exp(-m_j^2/M'^2)$ or $exp(-m_k^2/M'^2)exp(-m_j^2/M^2)$ and appears due to contraction of the poles associated with j_D or j_B. It has no exponential suppression in (M_D^2/M^2), in contradistinction with the normal situation (cf., e.g, ref.[5] or recent works in external fields [53,54]). This is the major shortcoming of our method in the problemat hand. Here we have to rely on suppression by powers of M_D^{-2} and numerical smallness.

Now we can represent the sum (37)-(39) in the following way

$$f_T(M^2, M'^2, Q^2) = K_1(M^2, M'^2, Q^2) + K_2(M^2, M'^2, Q^2) + K_3(M^2, M'^2, Q^2) \tag{40}$$

where the first term K_1 corresponds to the amplitude we are looking for, the second term K_2 corresponds to the contribution of the diagrams in which the quark currents j_D or j_B create intermediate havier than D or B; the third term, K_3, corresponds to the diagrams where the pole due to D or B mesons is absent. In other words, K_3 is generated by "wrong" cuttings (cuttings over intermediate states contaminating the sum rules for instance, $K\pi$ intermediate state in the j_D channel). More exactly,

$$K_1(M^2, M'^2, Q^2) = T e^{-m_D^2/M^2} e^{-m_A^2/M'^2} f_A f_B f_D \frac{m_D^2}{m_c} + \text{(higher resonance contribution in A channel)} \tag{41}$$

where T is the amplitude of $D \to AB$ decay.

$$K_2(M^2, M'^2, Q^2) = \sum_K \sum_n \sum_j T_{Knj} \frac{f_i f_n f_j}{(Q^2+m_j^2)} e^{-m_i^2/M^2} e^{-m_n^2/M'^2} +$$
$$+ \sum_{i,j,K,S} \frac{e^{-m_K^2/M'^2} e^{-m_i^2/M^2} T_{ij} T_{jKS}}{(Q^2+m_S^2)(m_i^2-m_j^2)} + \sum_{i,j,K,S} \frac{T_{iSj} T_{jK} f_i f_K f_S e^{-m_j^2/M^2} e^{-\frac{m_i^2}{M'^2}}}{(m_j^2-m_K^2)(Q^2+m_S^2)} \tag{42}$$

We will take into account the K_2 contribution in the framework of the standard continuum model[5, 57]. As for the function K_3, its appearance is most unpleasant and brings in the major theoretical uncertainty. It corresponds to the contribution of the diagrams where the pole connected with particles created by j_D or j_B, is contracted into the point. Somewhat symbolically,

$$K_3(M^2, M'^2, Q^2) = -\sum_j T_{Dj} T_{jAB} f_A f_B f_D \frac{m_D^2}{m_c} e^{-m_B^2/M'^2} e^{-m_j^2/M^2} -$$
$$- \sum_{i,j,K,S} \frac{T_{ij} T_{jKS} f_i f_K f_S e^{-m_j^2/M^2} e^{-m_K^2/M'^2}}{(Q^2+m_S^2)(m_i^2-m_j^2)} -$$
$$- \sum_j \frac{T_{DjA}}{Q^2} \frac{T_{jB} f_A f_B f_D m_D^2/m_c}{(m_K^2-m_j^2)} e^{-m_D^2/M^2} e^{-m_j^2/M'^2} - \sum_{i,j,K,S} \frac{T_{iSj} T_{jK} f_i f_K f_S e^{-\frac{m_j^2}{M'^2}-\frac{m_i^2}{M^2}}}{(Q^2+m_S^2)(m_K^2-m_j^2)} \tag{43}$$

It is easy to see that the second term in eq.(43) is strongly exponentially suppressed because the intermediate state j here contains heavy c-quark (thus, $m_j \geq m_D \gg m_B$). On the other hand, the first term in eq.(43) is suppressed much weaker than "normally": it contains a factor $\exp(-m_j^2/M^2)/(m_D^2-m_j^2)$ where m_j refers to light states instead of "normal" factor $\exp(-m_D^2/M^2)$ since here j contains only light quarks. It can be estimated numerically; the estimate will be given in our next work, here we will notice only that its contribution into the polarization operator seems to be less than 5% and in future we will take into account this term when estimating the uncertainty of the results obtained.

A few words about the channel connected with j_A . Here we do not apply the Borel transformation and must choose a certain model for the spectral density which would reflect the fact that, apart from π, K or η , there are higher resonance contamination. Apart from π , K or η we will account only for one next resonance, A_1, whith parameters which effectively include higher states as well. Qunatitatively, we will introduce in the r.h.s. of the sum rules the following formfactor $F_\pi(Q^2)$:

$$F_\pi(Q^2) = 1 - \frac{Q^2}{m_{A_1}^2+Q^2} \tag{44}$$

This model of the spectral density turns out to be self-consistent, as it will be seen below.

Transferring the continuum contribution to the l.h.s. we get the sum rule for the amplitude T:

$$f_T(M^2, M'^2, Q^2) = \frac{T f_D\, m_D^2/m_c\, f_A f_B\, F_\pi(Q^2)}{Q^2} e^{-(m_D^2/M^2 + m_B^2/M'^2)} \tag{45}$$

To reduce further the contribution to eq.(45) coming from the first term in eq.(43) we differentiate (45) with respect to $1/M^2$. Then the intermediate states with mass $m_j < m_D$ will be additionally suppressed by m_j^2/m_D^2. Finally, we arrived at

$$-\frac{1}{m_D^2}\frac{d}{d\,1/M^2}\;\frac{f_T(M^2,M'^2,Q^2)}{f_A f_B f_D\, m_D^2/m_c}\; e^{\left(\frac{m_D^2}{M^2}+\frac{m_B^2}{M'^2}\right)} = \frac{T F_\pi(Q^2)}{Q^2} \tag{46}$$

The full weak decay amplitude is T. To determine T we find the l.h.s. of eq.(46) for $Q^2 = 0.5 \div 1.6$ GeV^2 (at lower Q^2 higher power corrections blow up, at larger Q^2 higher resonances in the j_A channel may show up). The last stage of the procedure is more or less standard. We vary independently the Borel parameters M^2 and M'^2 and find a domain where the M^2, M'^2 dependence of the l.h.s. of eq.(46) is weak. At the same time we demand from the continumm contribution to be less than 30% and, simultaneously, higher powere corrections to be less than 30%. We check that Q^2 dependence isexhausted by $T F_\pi(Q^2)/Q^2$ and find the constant T, i.e. the $D \to AB$ decay amplitude.

4. Factorizable and Non-Factorizable Pieces in Weak-Decay Amplitudes

The sum rule (46) can be essentially simplified. Indeed, consider its theoretical part, $f_T(M^2,M'^2,Q^2)$. As has been already noted, for any two-particle decay it is the linear combination of three skeleton graphs, 3a,b, and c. Each of the graphs contains two blocks connected by the four-fermion interaction H_W. We can decompose f_T as follows:

$$f_T(M^2,M'^2,Q^2) = f_f(M^2,M'^2,Q^2) + f_n(M^2,M'^2,Q^2) \tag{47}$$

where $f_f(M^2,M'^2,Q^2)$ represents the sum of graphs with the two blocks connected only by H_W, with no gluon "bridge" from one block to the other.

The function f_n represents the sum of all other graphs.

Let us consider the issue in more detail: The function $f_f(M^2, M'^2, Q^2)$, as determined from Figs.3 a-c, is representable as a proruct of two terms fixed by separate blocks. In other words, we get the product of two correlation functions (a two-point function and a three-point function) calculated in the euclidean domain. Then we write down a dispersion relation for each of the correlation functions, saturate it by resonances, and, after the Borel transformation arrive (for example, for the function f_f determined by Fig.3b) at

$$f_f(M^2, M'^2, Q^2) = \frac{m_D^2}{Q^2} f_\pi f_+^K e^{-m_D^2/M^2} e^{-m_K^2/M'^2} +$$
$$+ \sum_i \frac{T_{\pi D i} \langle i | j_W | 0 \rangle \langle 0 | j_W | K \rangle}{m_i^2 - m_K^2} e^{-m_D^2/M^2} \left(e^{-m_i^2/M'^2} - e^{-m_K^2/M'^2} \right) + \tag{48}$$
$$+ (\text{higher resonance contribution})$$

where j_W denote axial vector (vector) currents comprising H_W. But this is nothing else than the sum rule for the quantity $m_D^2 f_\pi f_+^K$ i.e. for the factorized weak decay amplitude, plus the contribution of the diagram 5c, computed under the factorization hypothesis. (We omit here evident overall constants like $G_F C_1 \sin\theta$, etc). An analogous sum rule is valid for the graph 3c. Moreover, we get zero for the diagram 3c - manifestation of vanishing of the annihilation mechanism in the factorized amplitude (if $m_{u,d,s} = 0$). The arguments above referred to the channels $D, F \to PP$. They are equally applicable to the VP channels, however. Thus, the function $f_f(M^2, M'^2, Q^2)$ as it emerges from the sum rules, just reproduces the expectations of the standard model. Subtracting the sum rule for f_f from the total sum rule (46) we are left with the sum rule for the non-factorizable part of the amplitude giving deviations from the standard model

$$-\frac{1}{m_D^2}\frac{d}{d\frac{1}{M^2}}\frac{f_n(M^2,M'^2,Q^2)}{f_D m_D^2/m_c\, f_A f_B}\, e^{\frac{m_D^2}{M^2}+\frac{m_B^2}{M'^2}} = \frac{T_n f_\pi(Q^2)}{Q^2} \qquad (49)$$

This is our final result in this work.

To calculate the l.h.s. we keep only those graphs in which two blocks (Fig.3a-c) are connected by gluons. To the leading order the unit operator I and the operator $\bar{\psi}\psi$ are then irrelevant. The operators $\bar{\psi}(G\sigma)\psi$, G^2 , $\alpha_s(\bar{\psi}\psi)^2$ generated by graphs with gluons emitted inside one block are absorbed in $f_f^\rho(M^2,M'^2,Q^2)$. Hence, in our computations it is sufficient to keep dim 4,5 and 6 operators generated by graphs with gluon exchanges between the blocks.

5. Conclusions

Thus, we have constructed the sum rules for weak decay amplitudes of charmed mesons. Our final result, eq.(49), is formulated directly in terms of T_n - the non-factorizable piece giving the deviation from factorized prediction obtained within the spectator model. We have shown that in the theoretical part one must keep only the graphs with the gluon exchanges between the two blocks comprising the graph.

Unlike the previous models, all assumptions made can be checked within the approach itself, the error in the predictions can be estimated, and our accuracy is under the theoretical control. We expect $\Delta T/T \lesssim 20\%$. The main source of the uncertainty is the "other" resonance contamination.

Applications of the method to particular decays will be considered in the next paper.

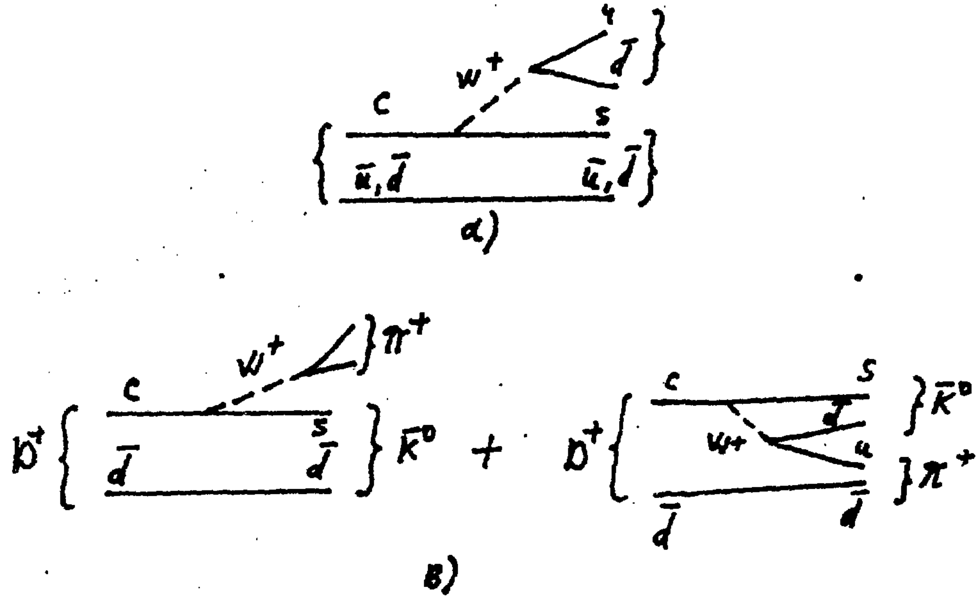

Fig.1. Spectator mechanism in the charmed-meson decays.

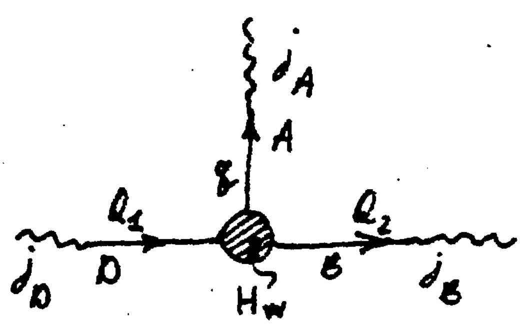

Fig.2. The four-point function induced by the quark currents and H_W. The shaded blob corresponds to the weak decay amplitude.

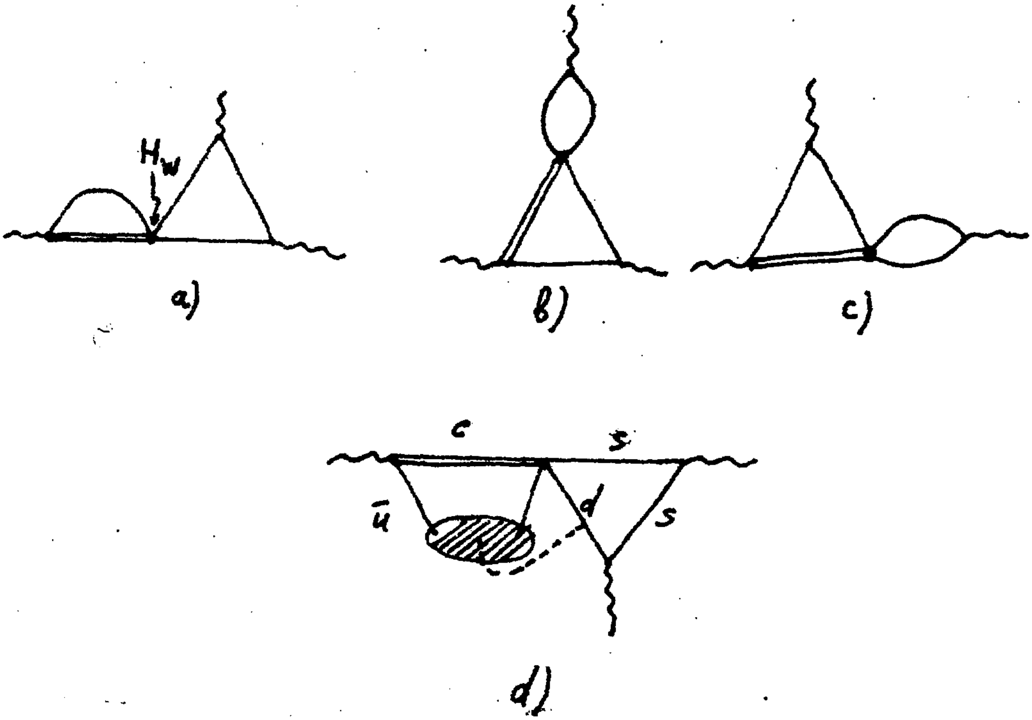

Fig.3. Theoretical part of the sum rules for charmed particle decays: (a-c) skeleton graphs determining the theoretical part; (d) an example of the "dressed" diagram for $D^0 \to \bar{K}^0 \phi$.

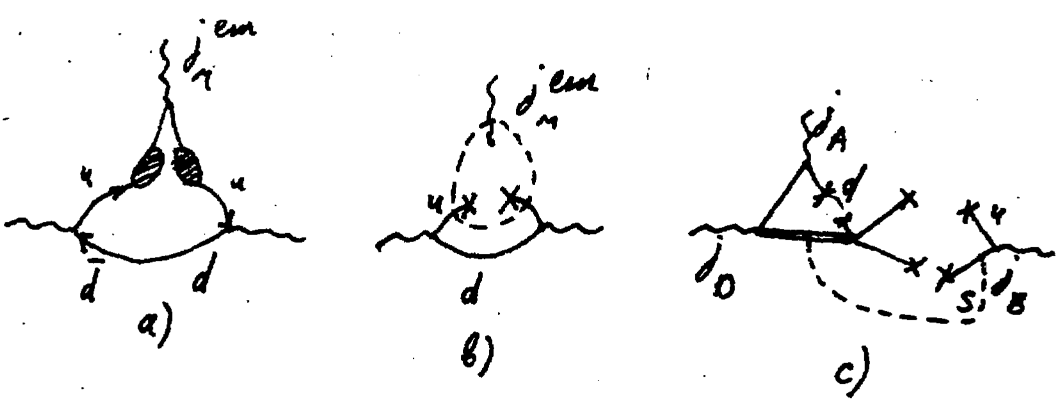

Fig.4. External-field-induced condensates.

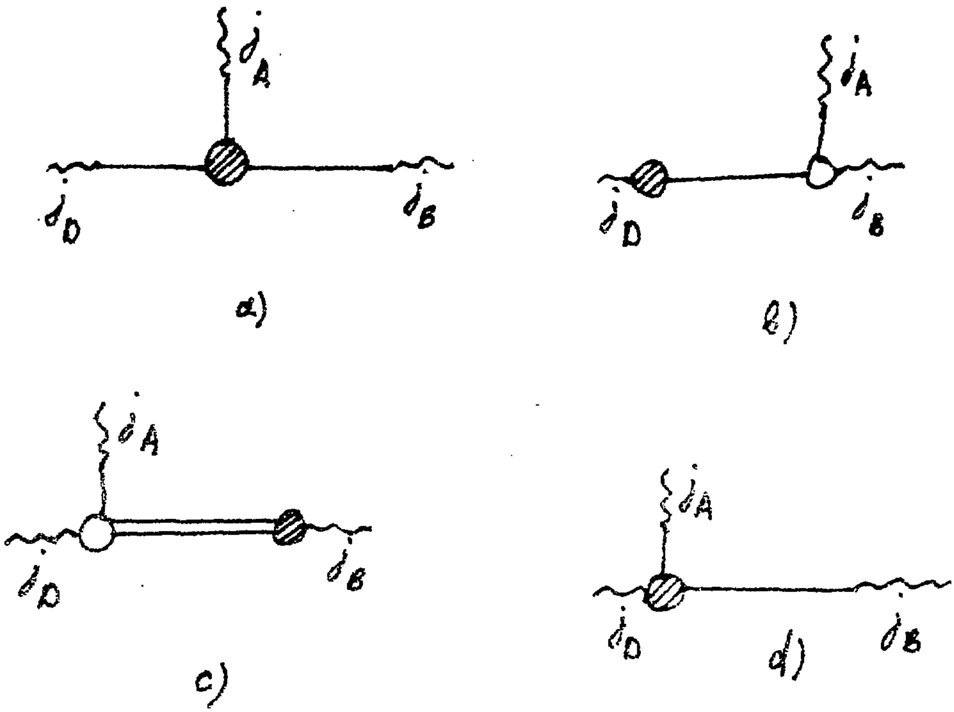

Fig.5. Phenomenological part of the sum rules for charmed meson decays. Here (and in Fig.7) we use the following notations: weak decay amplitude induced by H_W ; O strong decay amplitude.

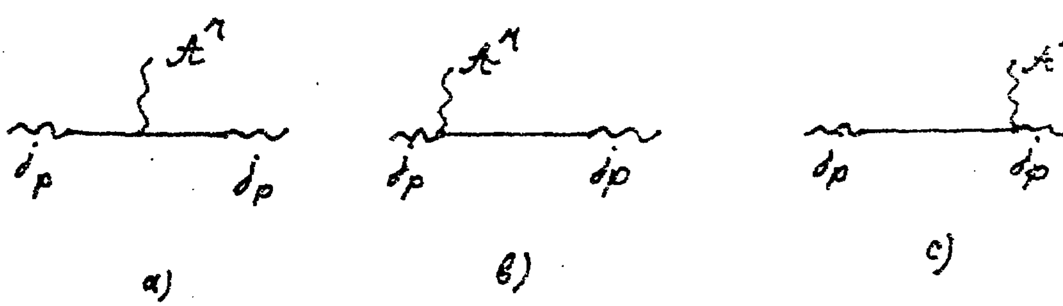

Fig.6. Phenomenological part of the sum rules for the proton-proton transition in the external axial field.

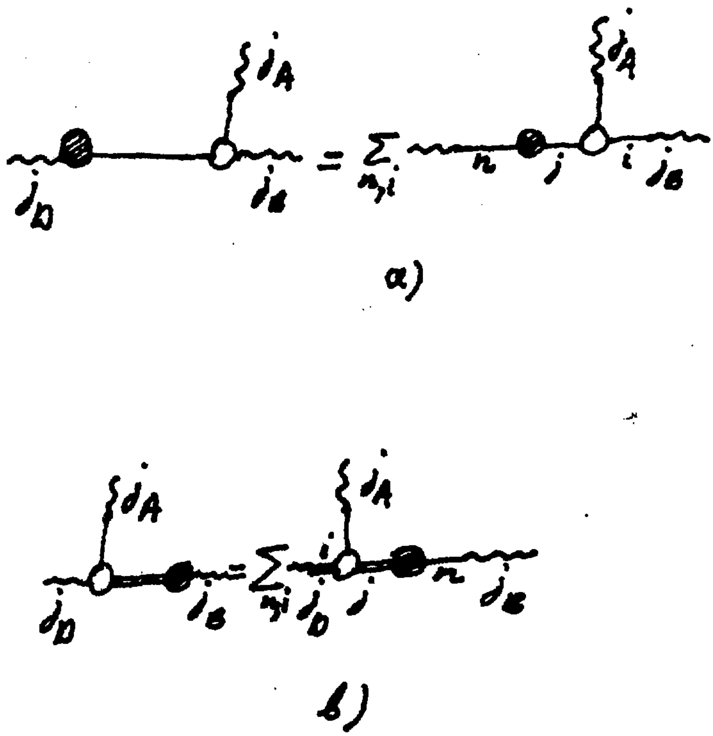

Fig.7. Pole model for the phenomenological part of the sum rules for weak decays of charmed mesons.

Weak decays of D and F mesons to be considered in the approach proposed.

	Cabibbo favoured decays	Cabibbo supressed decays
D^0	$\bar K^0\pi^+$, $\bar K^0\pi^0$, $\bar K^0\eta$, $\bar K^0\eta'$, $K^-\rho^+$, $\bar K^0\rho^0$, $K^*\pi^+$, $\bar K^{*0}\pi^0$, $\bar K^0\varphi$, $\bar K^0\omega$, $\bar K^{0*}\eta$	K^-K^+, $K^{-*}K^+$, K^-K^{+*}, $\pi^-\pi^+$, $\pi^0\pi^0$, $\rho^-\pi^+$, $\rho^0\pi^0$, $\bar K^0K^{0*}$, $K^0\bar K^0$, $\rho^+\pi^-$, $\pi^0\eta$, $\pi^0\eta'$, $\eta\Phi$, $\eta\omega$
D^+	$\bar K^0\pi^+$, $\bar K^0\rho^+$, $\bar K^{*0}\pi^+$	$\bar K^0K^+$, $\pi^+\pi^0$, $K^+\bar K^{0*}$, $K^{+*}\bar K^0$, $\pi^+\phi$, $\pi^+\omega$, $\pi^+\eta$, $\eta'\pi^+$, $\rho^+\pi^0$, $\pi^+\rho^0$
F^+	$\phi\pi^+$, $\eta\pi^+$, $\eta'\pi^+$, $\eta\rho^+$, $\bar K^0K^+$, $\bar K^{0*}K^+$, $\bar K^0K^{+*}$, $\omega\pi^+$, $\rho^0\pi^+$, $\pi^+\pi^0$	—

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Б.Ю.Блок, М.А.Шифман

К теории слабых адронных распадов очарованных мезонов.

Работа поступила в ОНТИ II.I2.85

Подписано к печати I7.I2.85 Т22385 Формат 60x90 I/I6
Офсетн.печ. Усл.-печ.л.2,0. Уч.-изд.л.I,4. Тираж 270 экз.
Заказ 9 Индекс 3624 Цена 2I коп.

Отпечатано в ИТЭФ, II7259, Москва, Б.Черемушкинская, 25

ИНДЕКС 3624

М., ПРЕПРИНТ ИТЭФ, 1986, № 9, с.1–32