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V.A.Khoze, M.A.Shifman,
N.G.Uraltsev, M.B.Voloshin

ON INCLUSIVE HADRONIC WIDTHS OF BEAUTIFUL PARTICLES

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V.A.Khoze, M.A.Shifman, N.G.Uraltsev, M.B.Voloshin - M.: ATOMIZDAT, 1986 - 16 p.

We consider preasymptotic (in the heavy quark mass) corrections in the non-leptonic decay rates of beautiful mesons and baryons. The effects of real and virtual gluon emission are taken into account. Main emphasis is put on the difference of the lifetimes of B_s^0 mesons with different CP parity. This difference is found to be the largest in the b-family. Possibilities for experimental study of this phenomenon are briefly discussed.

Fig. - 8, ref. - 16

I. Introduction

The recent paper[1] presents a systematic analysis of preasymptotic effects in the inclusive weak hadronic decays of charmed and beautiful particles. (The history of the problem and a list of references can be found, e.g., in the review paper[2]). The results of ref.[1] can be summarized as follows.

(i) In the charmed particle family the preasymptotic effects are of order unity. The sign of the corrections is determined reliably while their magnitude can be estimated semiquantitatively. Thus, the hierarchy of lifetimes of D^+, D^0, F, Λ_c^+, $\Xi_c^{+(A)}$, Ω_c^0 is predicted (the particles in this list are ordered according to the expected lifetimes; D^+ is the most shortliving).

(ii) For the beautiful hadrons the difference in decay widths is found to be much smaller. A conservative bound[1] is:

$$\Delta\Gamma(\text{b hadrons})/\Gamma \lesssim 0.05 .$$

In this paper we add a few points not covered by the analysis of ref.[1]. These points concern a hard gluon emission, the mechanism proposed by Bander et al.[3]. We dwell also on the $B^0\bar{B}^0$ mixing which results in a sizeable effect for the B_s^0 mesons (see e.g. ref.[4]). It will be shown that the graphs with emission of a hard gluon[3] are rather inessential for the lifetimes of the charmed hadrons; in the b-hadron family, however, their effect can be comparable to that of other corrections and can shift the widths by O(1%). Therefore the conclusion[1] about small difference of lifetimes of the b hadrons is still valid, while the precise ordering of the total rates given in ref.[1] (see eq.(6) in[1]) can be somewhat modified. As to the mixing of the neutral B mesons, the effect can lead to the largest lifetime difference, $\Gamma(B_s^0$ "shortliving") − − $\Gamma(B_s^0$ "lonliving"), the largest among all beautiful hadrons). $\Delta\Gamma$ in this case can amount to $\sim 0.1\,\Gamma(B_s^0)$.

2. $B-\bar{B}$ oscillations and the difference of the decay widths

As well known, in the limit of large quark mass $m_Q \to \infty$ the non-leptonic decay rate is determined by the free quark decay (the so called spectator graph, Fig.1). In this approximation the lifetimes of all beautiful hadrons are the same. Non-universality of the widths arises due to various preasymptotic effects induced by the presence of the light quark. For example, in the B^+ decay one should take into account the interference of the u quark produced in the decay $\bar{b} \to \bar{c} u \bar{d}$ with the spectator u quark (Fig.2). For the neutral B meson states like $\bar{b}d$ or $\bar{b}s$ this mechanism is not applicable directly and the preasymptotic corrections are generated by annihilation graphs. The latter have not been discussed in ref.[1] and we will dwell on the issue in Sec.3. Here let us notice the following. Due to $B-\bar{B}$ transitions the eigenstates are not $\bar{b}s$ and $b\bar{s}$, but, rather, their linear combinations

$$B_+ = \frac{1}{\sqrt{2}}(\bar{b}s - b\bar{s}) \quad ; \; CP = +1, \qquad (1)$$

$$B_- = \frac{1}{\sqrt{2}}(\bar{b}s + b\bar{s}) \quad ; \; CP = -1 .$$

(Only B_s^0 mesons are considered here since only for them the mixing is essential, $\Delta m \gtrsim \Gamma$, while for B_d^0 $\Delta m \ll \Gamma$ is expected. We also neglect small CP violation). Since the wave function of $B_+(B_-)$ is a coherent mixture of $\bar{b}s$ and $b\bar{s}$ the interference of the s ($\bar{s}$) quarks appears in the decay (Fig.3). As a result the widths of B_+ and B_- are shifted in the opposite direction, and their difference is given by

$$\Delta\Gamma \equiv \Gamma_{non\text{-}lept.}(B_+) - \Gamma_{non\text{-}lept.}(B_-) \simeq -2\,\frac{Re(\Gamma_{12} M_{12}^*)}{|M_{12}|} \quad (2)$$

where Γ_{12} and M_{12} are the absorptive and the dispersive parts of the mass matrix element [4,5]. If CP violation is neglected then

$$\Delta\Gamma = -2\Gamma_{12} \quad . \quad (2a)$$

(A general treatment of the two-level system can be found e.g. in ref. [5]). The dominant contribution to Γ_{12} is given by the imaginary part of the graphs of Fig.3. It can be readily noted that the contribution of the graph 3b is suppressed by chirality factor $\sim m_c^2/m_b^2$. Calculation of the effect is quite straightforward (see papers [6-9] and references therein). Neglecting for a while gluon corrections one finds

$$\Gamma_{12} \simeq \frac{G_F^2}{4M_B}(V_{cb}V_{cs}^*)^2\, T_{\mu\nu} \langle \bar{B}_s^0 | \bar{b}_i \Gamma_\mu s_j\ \bar{b}_j \Gamma_\nu s_i | B_s^0 \rangle \quad (3)$$

where $\Gamma_\mu = \gamma_\mu(1+\gamma_5)$, M_B is the B meson mass, V_{cb}, V_{cs} are the Kobayashi-Maskawa parameters, and $T_{\mu\nu}$ is the absorptive part of vacuum polarization by the current $\bar{c}\Gamma_\mu c$ at $p^2 \simeq m_b^2 \simeq M_B^2$:

$$T_{\mu\nu} = \frac{1}{6\pi}\left(1-\frac{4m_c^2}{p^2}\right)^{1/2}\left((p^2-m_c^2)g_{\mu\nu} - \left(1+\frac{2m_c^2}{p^2}\right)p_\mu p_\nu\right) \quad . \quad (4)$$

Here p_μ is the momentum of the meson (or of the b quark). Using the standard factorization procedure for the matrix element one finds

$$\Gamma_{12} = \frac{G_F^2}{8\pi} M_B^3 f_B^2 \left(V_{cb} V_{cs}^* \right)^2 \tag{5}$$

where m_c is neglected and f_B is the axial constant of the B_s^0:

$$\langle 0 | \bar{b} \gamma_\mu \gamma_5 s | B_s^0 \rangle = i f_B p_\mu$$

Eq.(5) corresponds to vanishing mass of the charmed quark. The correction factor due to m_c amounts to $(1-4m_c^2/M_B^2)^{1/2} \times [1 - \frac{2}{N_c} \frac{m_c^2}{M_B^2}] \approx 0.75$.

Gluon exchnages give rise to a renormalization factor in eq.(3), which is traditionally denoted by η. We consider here this factor only in the leading logarithmic approximation (LLA). There are two types of such corrections. The domain of momenta of virtual gluons $m_b^2 < p^2 < m_W^2$ gives rise to the factors $C_\pm$ in the $\Delta B = 1$ weak Lagrangian [10]

$$C_- = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_W)} \right)^{12/23} \approx 1.29; \; C_+ = (C_-)^{-1/2} \approx 0.88 .$$

The corresponding effect in Γ_{12} is well known (see e.g. refs. [6,7]) and requires no further comments. Less known are the so-called hybrid logs [11] which come from the domain of momenta $\mu^2 < p^2 < m_b^2$ where μ is the characteristic quark off-shellness in the B meson ($\mu \sim$ 300 MeV). Details about the hybrid logs can be found in refs. [1,11]. The hybrid renormalization of the operator $O = (\bar{b}\Gamma_\mu s)(\bar{b}\Gamma_\mu s)$ was calculated in [11]. It was found that this operator is multiplicatively renormalized in LLA and its hybrid anomalous dimension, $\gamma_O = 4$, is exactly twice that of the current $j_{\mu 5} = \bar{b}\gamma_\mu\gamma_5 s$ ($\gamma_{j_{\mu5}} = 2$).

Since factorization for the four-quark matrix element is expected to be valid at a low normalization point ($\sim\mu$), the evolution of the operator from the normalization point m_b down to μ should be explicitly taken into account. Notice that the constant f_B in eq.(5) is defined for the current normalized at a high point ($\sim m_W$) (or, which is equivalent in LLA, $\sim m_b$). This fact and the above mentioned relation $\gamma_0 = 2\gamma_{j_{\mu 5}}$ leads to the conclusion [11] that the matrix element $\langle \bar{B}_s^0 | (\bar{b}\Gamma_\mu s)^2 | B_s^0 \rangle$ when expressed in terms of f_B, contains no hybrid logs.

However, this is not the whole story for the matrix element in eq.(3) since $T_{\mu\nu}$ contains also a term $p_\mu p_\nu$, and one has to consider the evolution of the operator

$$O_1 = (\bar{b}\Gamma_\mu s)(\bar{b}\Gamma_\nu s). \tag{6}$$

Since the hybrid logs arise only due to the color-charge interaction of the heavy quark [11] O_1 mixes only with the following operator:

$$O_2 = (\bar{b}_i \Gamma_\mu s_j)(\bar{b}_j \Gamma_\nu s_i) \tag{6a}$$

In LLA the combinations diagonal with respect to renormalizations are:

$$O_+ = \frac{1}{2}(O_1 + O_2), \quad \gamma_+ = \gamma_0 = 3\frac{N_c^2 - 1}{2N_c} = 4, \tag{7}$$

$$O_- = \frac{1}{2}(O_1 - O_2), \quad \gamma_- = \gamma_0 - 2\frac{N_c + 1}{N_c} = \frac{4}{3}.$$

Therefore, the effect of the hybrid log renormalization reduces to the substitution

$$O_{1\mu\nu} \to \frac{1}{2}\varkappa^{\gamma_+/b}\Big[(1+\varkappa^{\Delta\gamma/b})O_{1\mu\nu} + $$
$$+ (1-\varkappa^{\Delta\gamma/b})O_{2\mu\nu}\Big],$$
$$O_{2\mu\nu} \to \frac{1}{2}\varkappa^{\gamma_+/b}\Big[(1-\varkappa^{\Delta\gamma/b})O_{1\mu\nu} \qquad (8)$$
$$+ (1+\varkappa^{\Delta\gamma/b})O_{2\mu\nu}\Big],$$
$$O \to \varkappa^{\gamma_+/b} O\ ;\ \varkappa = \alpha_s(\mu)/\alpha_s(m_b),\ \gamma_+ = 4;\ \Delta\gamma = \gamma_- - \gamma_+ = -\frac{8}{3}.$$

Here b is the first coefficient of the QCD β -function, $b \approx 25/3$ (we neglect the small step in b at the charm threshold). Using the factorization procedure and the Fierz transformation

$$(\bar\psi_1\Gamma_\mu\psi_2)(\bar\psi_3\Gamma_\nu\psi_4) = -\frac{1}{2}\Big\{\bar\psi_1\Gamma_\mu\psi_4\ \bar\psi_3\Gamma_\nu\psi_2 +$$
$$+ \bar\psi_1\Gamma_\nu\psi_4\ \bar\psi_3\Gamma_\mu\psi_2 - g_{\mu\nu}\ \bar\psi_1\Gamma_\alpha\psi_4\cdot\bar\psi_3\Gamma_\alpha\psi_2 \qquad (9)$$
$$- i\varepsilon_{\mu\nu\rho\lambda}\ \bar\psi_1\Gamma_\rho\psi_4\ \ \bar\psi_3\Gamma_\lambda\psi_2\Big\}$$

we find in the limit of vanishing c quark mass

$$\Gamma_{tz} = \frac{G_F^2}{8\pi} M_B^3 f_B^2\, \eta\, (V_{cb}V_{cs}^*)^2,$$
$$\eta = \Big\{\frac{1}{2}(C_+^2 + C_-^2) + \frac{1}{2N_c}(C_+^2 - C_-^2) + \frac{1}{2}\Big(1-\frac{1}{N_c}\Big)\times \qquad (10)$$
$$\times(1-\varkappa^{\Delta\gamma/b})\Big[N_c\Big(\frac{C_+ - C_-}{2}\Big)^2 + \frac{1}{2}(C_+^2 - C_-^2) -$$
$$-\Big(\frac{C_+ + C_-}{2}\Big)^2\Big]\Big\} \approx 0.86 \qquad \text{(for } N_c=3)$$

The suppression factor due to the c quark mass is ≈ 0.8. Thus the gluon corrections slightly suppress the bare estimate (5). For a numerical estimate it is convenient to normalize $\Delta\Gamma = 2\Gamma_{12}$ with respect to the semileptonic rate

$$\Gamma_{se} \equiv \Gamma(b \to \mu\nu_\mu c) = \frac{G_F^2 m_b^5}{192\pi^3} |V_{cb}|^2 Z_c ,$$

$$m_b^5 Z_c \simeq (4.8\ \mathrm{GeV})^5 . 0.55$$

(this estimate of the phase-space factor Z_c times m_b^5 is close to the one found in a specific model [12] from the analysis of experimental lepton spectra [13], $m_b^5 Z_c = (5\ \mathrm{GeV})^5 . 0.4$). Thus, for the ratio $\Delta\Gamma/\Gamma_{se}$ we obtain the estimate

$$\frac{\Delta\Gamma}{\Gamma_{se}} \approx 48\pi^2 \frac{f_B^2}{M_B^2} \cdot 2{,}0 . \qquad (12)$$

For the parameter f_B we use the value $f_{B_s} \approx 130$ MeV which is somewhat larger than the average theoretical estimate [14,15] for f_{B_d} (see also ref. [2]). Since experimentally $B_{se} = \Gamma_{se}/\Gamma_{tot} \approx 0.12$ we finally obtain the estimate

$$\Delta\Gamma/\Gamma_{tot} \approx 0.07 . \qquad (13)$$

A general remark is in order here. The logarithmic renormalization of Γ_{12} is discussed here just in the same way as that of M_{12} in ref. [11]. Namely, the absorptive part of the quark diagram is treated as a local operator whose renormalization is considered in a standard way. Naturally, one may wonder whether this procedure is adequate for the absorptive part, at least in LLA.

For example, the graphs of Fig.4 are obviously ignored in this approach. The graph 4a corresponds to emission of a real gluon while that in Fig.4b describes the vertex correction due to a virtual gluon. Each of these diagrams contains an additional logarithm. This logarithmic factor, however, cancells in the sum of the graphs. Such cancellation between real and virtual gluons along each quark line in LLA is a general phenomenon *) well-known to the experts in hard QCD. We have also checked the validity of the procedure used in a simplified model.

3. Gluonic mechanism in decays of heavy mesons.

In ref.[3] an annihilation mechanism in inclusive D^0 decays have been discussed. The bare quark diagram is shown in Fig.5a. The corresponding contribution to the width is suppressed with respect to the free quark decay by the factor

$$\Gamma_{ann}/\Gamma_{se} \approx \frac{1}{N_c} 24\pi^2 \frac{f_D^2}{M_D^2} \frac{m_s^2}{M_D^2} \rho(C_+, C_-) . \qquad (14)$$

Here f_D^2/M_D^2 reflects the fact that in theannihilation process c and $\bar{u}$ quarks should be at the same point, and the factor m_s^2/M_D^2 appears due to the chirality suppression. As has been noticed[3], the emission of a gluon (Fig.5b) eliminates the chirality suppression factor. Furthermore, emission of a gluon by the light quark replaces the factor f_D^2/M_D^2 by f_D^2/E^2, where E is the characteristic energy of the light quark inside the meson, $E \sim 300$ MeV. The rea-

*) We acknowledge useful discussions of this and related points with Yu.L.Dokshitzer.

son for such an enhancement is that the bare quark annihilation (Fig.5a) occurs in the volume $V \sim M_D^{-3}$ while the off-shellness of the light quark in the graph of Fig.5b is $\sim m_c E \ll m_c^2$; hence, the decay proceeds in a larger volume. The contribution of the mechanism of Fig.5b is approximately given by[3]

$$\frac{\Delta\Gamma}{\Gamma_{se}} \approx \frac{8\pi}{27}\alpha_s \frac{f_{D,B}^2}{E^2}\left(\frac{C_+ + C_-}{2}\right)^2 \simeq \begin{cases} 0.14 & \text{for } D \\ 0.05 & \text{for } B \end{cases} \quad (15)$$

Clearly, the effect is numerically small for D mesons in comparison with the corrections discussed in ref.[1]. Though parametrically the ratio (15) is $O(m_Q^{-1})$, and the corrections of ref.[1] are $O(m_Q^{-3})$, there is a small numerical factor in eq.(15) which reflects the three-body phase space suppression with respect to the two-body one involved in the corrections of ref.[1].

For beautiful hadrons the both effects can be of comparable strength. This remark refers not only to the mechanism of Fig.5b*) but in general to all effects of real gluon emission (we mean hard gluons). These effects do not give log factors, but are enhanced by the factor M/E in the amplitude. Two typical examples of such effects are shown in Fig.6.

Fig.6a displays one of the graphs describing the interference correction to the D^+ decay rate. If the gluon line is removed the graph turns into the one considered in ref.[1] which gives (up to log renormalization effects)

$$\Delta\Gamma(D^+)/\Gamma_{se} \approx -N_c^{-1}\, 48\pi^2 f_D^2 M_D^{-2} .$$

*) In annihilation contributions to the decay widths of charmed mesons the main role is played by non-perturbative gluons. According to ref.[16], the annihilation mechanism gives ~ 20% enhancement of the widths of D^0 and F.

The gluon in Fig.6a does not produce log factor, but due to the pseudoeclidean kinematics it produces M/E factor which removes, to a certain extent, the suppression factor f_D^2/M_D^2 . As a result, the ratio of $\Delta\Gamma/\Gamma_{sl}$ given by the graph 6a to the one calculated in ref.[1] is of the order of

$$\frac{N_c^2-1}{N_c}\ \frac{\alpha_s}{16\pi}\left(\frac{M_D}{E}\right)^2 \approx 2\alpha_s < 1 \qquad (16)$$

(Here α_s is normalized at $\mu\sim(Em_c)^{1/2}$).

Thus the mechanism[3] brings some uncertainty in the estimates of the preasymptotic corrections to the decay rates, which is, however, smaller than the expected overall uncertainty of the estimates of ref.[1] for the charmed particles. Therefore the main conclusions of ref.[1] are not altered by this mechanism.

For beautiful hadrons the situation is somewhat different and the gluonic diagrams can be of the same order and even exceed the skeleton ones. E.g., the contribution of a graph similar to Fig.6a to the width of B^+ is approximately given by

$$-\left(1-\frac{1}{N_c^2}\right)\pi\alpha_s\ \frac{f_B^2}{E^2}\ \Gamma_{sl}^0 \sim -0.4\alpha_s\Gamma_{sl}^0 \qquad (17)$$

which is about 1% of the total rate. Therefore the hard gluon emission can considerably affect the quantitative estimates of ref[1] for b hadrons. However, the conclusion[1] that the spread of the decay rates of the b-hadrons is $\lesssim 5\%$ seems to be still valid. Notice also that we have considered only the graph of Fig.6a. There are, however, other diagrams describing the same mechanism (see e.g.,Fig.7), and some cancellation between these graphs is possible.

It is also interesting to notice that though the annihilation mechanism of Fig.5a and the interference one of Fig.6a are parametrically of the same order the former one is numerically suppressed. As to the hard gluon correction (non-logarithmic) to

$\Delta\Gamma$ of B_+ and B_- (see the skeleton graph of Fig.3) it has been calculated in a non-relativistic model in ref.[7]. It has been found that the contribution $O(M_B^2/E^2)$ completely cancells out (this can be understood by considering the gluon polarization in the transitions of B and $\bar{B}$ into $c\bar{c} + g$). Therefore, the ratio of the diagram 8 to that of Fig.3 is of order

$$\frac{\alpha_s}{16\pi}\ \frac{M_B}{E} \lesssim \quad 10\% \quad . \qquad (18)$$

This cancellation might be eliminated by emission of an additional gluon; however, the extra factor α_s and the smallness of the phase space for the $c\bar{c} + 2g$ transition must strongly suppress this contribution. Thus, it is likely that the difference of the widths of B_+ and B_- is to a good accuracy given by eqs.(10),(11).

4. Conclusions

The problem of preasymptotic effects in decays of heavy hadrons is important for hadronic physics. Confronting theoretical estimates with experimental data one can test various models of quark dynamics in heavy hadrons and in the weak decay. As a convincing example one can recall the theoretical efforts stimulated by the unexpected discovery of the large difference in D^0 and D^+ lifetimes.

In contrast with the charmed particles the weak decay widths of the beautiful hadrons should be almost equal to each other; the maximal difference is expected for B_s^0 mesons. Therefore, we briefly discuss here some possibilities of observing the difference $\Gamma_+ - \Gamma_-$. First of all it can be noticed that all standard (integral) characteristics of the B-$\bar{B}$ system depend on the square of the parameter $\Delta\Gamma/\Gamma$ (see e.g. ref.[4]) which makes it practically impossible to determine $\Delta\Gamma$ from measurements of such parameters. The value of $\Delta\Gamma$ figures also in the time dependences of different final states in B decays. (The time is determined by the decay length). Unfortunately, in this case the effects due to $\Delta\Gamma$ also turn out to be quadratic in $\Delta\Gamma/\Gamma$ since it is practically impossible to trace the B decays at time intervals considerably larger than Γ^{-1}. For example, the time dependence of the lepton yield in the tagged B beam is as follows

$$N\left(B^0 \to \ell^{\pm}(t)\right) \sim \frac{1}{2} \exp(-\Gamma t)\left[\cosh\frac{\Delta\Gamma t}{2} \pm \cos\Delta M t\right],$$

$$\Delta\Gamma = \Gamma_+ - \Gamma_- , \quad \Gamma = \frac{1}{2}\left(\Gamma_+ + \Gamma_-\right). \tag{19}$$

Linear in $\Delta\Gamma$ effects appear only if certain states f with definite CP parity are studied. Assuming CP conservation, B decays into such states are described by the standard exponential factors $\exp(-\Gamma_+ t)$ and $\exp(-\Gamma_- t)$. Naturally, the measurement of the difference between these exponents is most reliable when the states of different CP parity are detected in one and the same experiment. Such a situation could be realized in a correlation-type experiment. For example, in the process $e^+e^- \to \gamma^* \to B^0\overline{B^0}$ only the combination B_+B_- is produced. Therefore, one can tag one of the mesons by detecting a state f_+ or f_- with the definite CP and measure the time dependence, say, of the lepton yield in the decay of the second B meson.

Another possibility arises due to CP nonconservation in the decays $B^0 \to f$. It can be realized by measuring the time dependence of production of a CP eigenstate f in a tagged B_s^0 beam [4].

This dependence is approximately given by

$$N\left(B_0\to f(t)\right)\sim \exp(-\Gamma t)\left[\cosh\frac{\Delta\Gamma t}{2}+\right.$$
$$\left.+\,\mathrm{Re}\,\lambda\,\sinh\frac{\Delta\Gamma t}{2}-\mathrm{Im}\,\lambda\,\sin(\Delta M t)\right] \qquad (20)$$

where $\lambda = -\nu e^{2i\phi}$, ν is the CP parity of the state f , and ϕ is the CP odd phase of the amplitude $\bar{B}^0\to f$ which for some decays can be quite sizeable [4]. As it is seen from eq.(20) the sign of the effect produced by $\Delta\Gamma$ is correlated with the CP parity of f.

To summarize, we have made a more accurate estimate of $\Delta\Gamma$ for B_s^0 mesons with definite CP-parity. In this system the difference of lifetimes seems to be maximal in comparison with other weakly-decaying beautiful hadrons. Apart from that we have discussed the effects of hard (non-logarithmic) gluons [3] in various preasymptotic corrections to the decay rates. For charmed particles these effects are found to be small, while for beautiful hadrons hard gluons can somewhat modify the quantitative conclusions of the previous work [1]. However, the conclusion [1] that the difference in the weak decay widths of b hadrons is not larger than a few per cent, remains intact.

We thank Ya.I.Azimov and Yu.L.Dokshitzer for valuable discussions.

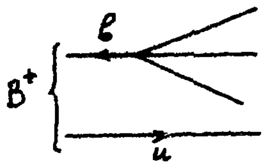


FIG.1

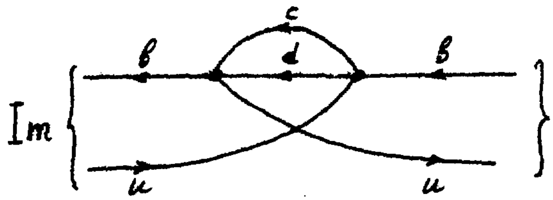
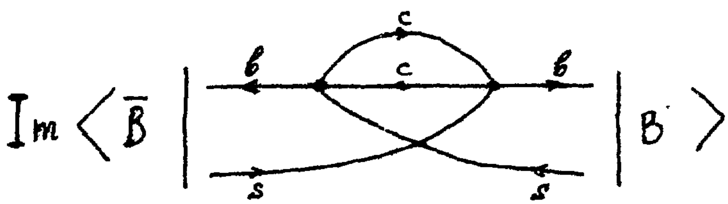


FIG.2



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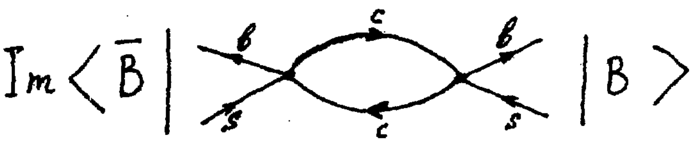
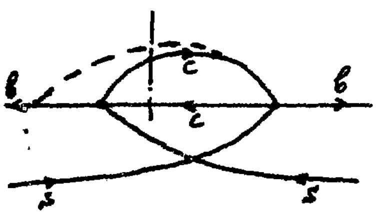
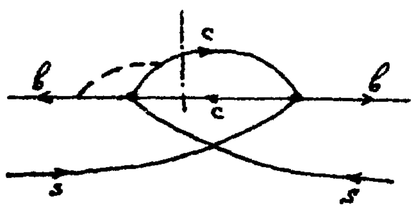


b

FIG.3



a



b

FIG.4

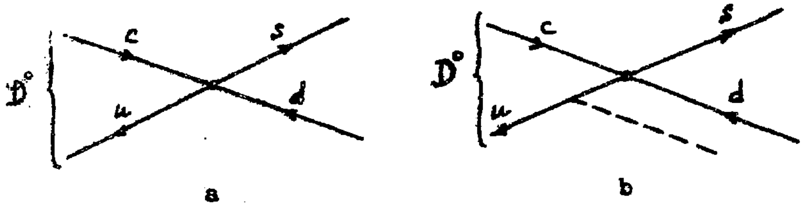


FIG.5

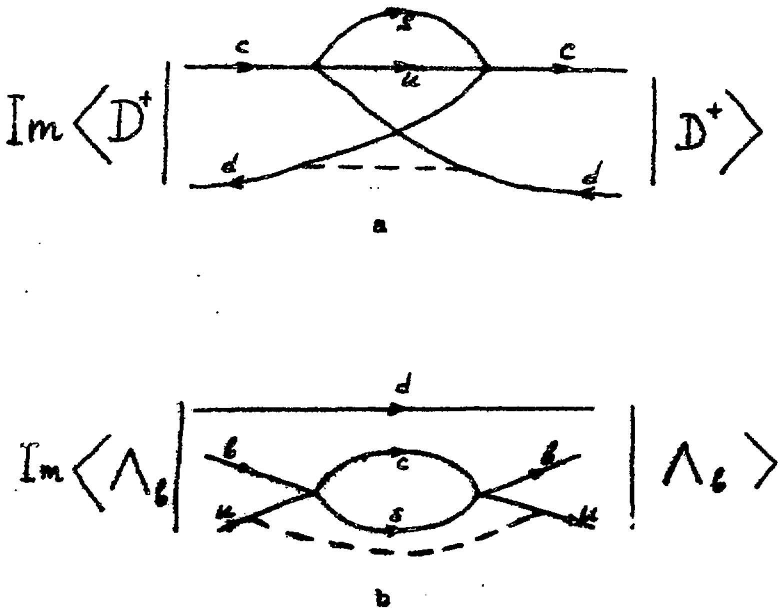


FIG.6

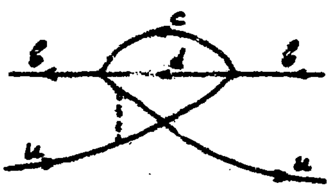


FIG.7

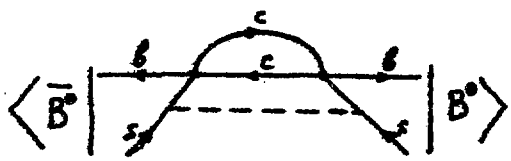


FIG.8

References

1. Shifman M.A., Voloshin M.B. **Moscow, Preprint ITEP,1986, N 83.**

2. Shifman M.A., Talk at the Int.Symp.on Production and Decay of Heavy Hadrons, Heidelberg, 1986.

3. Bander M., Silverman D., Soni A.,-Phys.Rev.Lett.,1980,44, 7.

4. Azimov Ya.I., Uraltsev N.G., Khoze V.A., Preprint LNPI-11 88, 1986.

5. Lee T.D., Wu C.S.-Ann.Rev.Nucl.Sci.,1965,15,381; 1966,16,471.

6. Hagelin J.S.-Nucl.Phys.,1981,B193, 123.

7. Franco E., Lusignoli M., Pugliese A.-Nucl.Phys.1982,B194,403.

8. Chau L.L., Keung W.Y., Tran M.D.-Phys.Rev. 1983,B27,2776.

9. Chau L.L.-Phys.Reps. 1983, 95, 1.

10. Gaillard M.K., Lee B.W.-Phys.Rev.Lett. 1974, 33,138; Altarelli G., Maiani L.-Phys.Lett., 1974,52B,354.

11. Shifman M.A., Voloshin M.B. **Moscow, Preprint ITEP, 1986, N 54.**

12. Altarelli G., Cabibbo N., Corbro G. et al.- Nucl.Phys.1982, B238, 365.

13. Klanner R. Talk given at the XXII Int.Conf. on High Energy Physics, Leipzig, July 19-25, 1984, v.1, p.201.

14. Zhitnitsky A.R., Zhitnitsky I.R., Chernyak V.L.-Yad.Fiz.,1983, 38, 1277.

15. Aliev T.M., Eletsky V.L.-Yad.Fiz., 38, 1537.

16. Blok B.Y., Shifman M.A. **Moscow, Preprint ITEP,1986,N 9,17,37**

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