Ref.TH.2929-CERN

CP VIOLATION IN DECAY RATES OF CHARGED BOTTOM MESONS

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A B S T R A C T

In the standard theory we discuss a mechanism of interference between two different tree-level charged current amplitudes to induce CP violation. The asymmetry between conjugate Cabibbo disfavoured modes of charged bottom meson decays is considered. Estimates for two-body decays of B_u and B_c are presented. We find that $B_u^- \to D^- + D_o^*$, for instance, can give a big CP asymmetry with values $\sim$1-50%, depending on the angle and phase parameters.

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Ref.TH.2929-CERN
19 August 1980

1. - INTRODUCTION

Ever since the discovery [1] of CP violation in 1964, the effect has only been seen [2] in the K^o-$\bar{K}^o$ system and its origin is not understood. At present, the most attractive model of CP violation is that of Kobayashi-Maskawa [3] (hereafter denoted by KM) where the existence of three families of quarks allows the presence of a CP violating phase in the charged current sector, together with three Cabibbo-like rotation angles. With the prejudice that all rotation angles should be small, the predictions [4] of the KM model for the K^o-$\bar{K}^o$ system are rather similar to those of the superweak model [5]. A small departure from the superweak prediction $\varepsilon' = 0$ is, however, anticipated and one hopes that the precision experiment [6] now in progress at Fermilab will provide a test of the KM model.

In connection with heavy quarks, a great deal of theoretical effort has been spent [7] on studying the CP properties of the B^o-$\bar{B}^o$ system. In view of the great importance of understanding the origin of CP violation one would like to study situations in which the effect may exhibit itself directly in a transition. Therefore, in this paper, we address ourselves to the question of CP violation in the decay of charged bottom mesons. A particularly clean signature of CP violation would be [8] the demonstration of inequality of the decay rate or branching ratio for any decay process and the corresponding quantity for its charge-conjugated process. The question is whether one expects such a difference and if so at what level in the standard theory.

We show that Cabibbo-disfavoured modes of B meson decays may present big CP asymmetries, as a consequence of having two interfering amplitudes of similar strength with different weak interaction phases. In Section 2, general considerations about CP effects in the standard theory are made. It appears that tree-level CP violating amplitudes are only possible from heavy quark systems, and the pieces of the weak Hamiltonian responsible for the transitions are given. Section 3 explores the effect in some two-body decays of B_u and B_c pseudoscalar mesons, using a constituent quark model for the meson form factors. The estimates are given in Section 4, as a function of the KM parameters, for $B_u^- \to D^- + D_o^*$ and $B_c^- \to \pi^- + \bar{D}_o$. Section 5 summarizes the main points and presents some conclusions.

2. - CP ASYMMETRIES

A glance at the KM matrix tells us that the dominant transitions *) for B mesons are from the $b \to c$ vertex for which the imaginary part of the transition coupling constant is, a priori, as important as the real part. Thus one may wonder why CP violation in B decays is not enormous. There are, however, two points.

The first point is that an over-all phase is not what matters but relative phases between two (or more) amplitudes contributing to the same process. As all other interactions, strong, electromagnetic and neutral currents, are flavour-conserving, the only possibility to build up a relative phase is to have interference with another charged current flavour changing amplitude. This implies that, for weak interactions operating to the lowest order and at the tree level, one will have CP violating effects mostly in the Cabibbo suppressed decays and not in the allowed decays. This is similar to what happens in atomic and nuclear physics where one looks for parity violating effects in "forbidden" transitions. If instead one starts with a Cabibbo allowed decay then the corresponding CP violating amplitude is doubly forbidden as compared to the allowed one. So we study the inferference between two different tree-level weak interaction amplitudes. A different mechanism proposed [9] involves strong radiative corrections to a loop giving flavour changing neutral currents.

The second point is that even if two interfering weak interaction amplitudes have different phases the rate for a process and its charge conjugated process will be equal by CPT as long as the weak transition matrix is hermitian (i.e., lowest order) and the strong interactions are neglected, viz.,

$$\langle f | T | i \rangle = \langle \bar{f}' | T | \bar{i}' \rangle^* \tag{1}$$

where $\bar{i}'$ is the antiparticle of i with its spin reversed. In fact, in order to get CP violation in the rates it is necessary to have at least two interfering amplitudes with different weak interaction <u>as well as</u> different strong interaction phases. By writing the T matrix as

$$T = g_1 T_1 + g_2 T_2 \tag{2}$$

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*) We refer to $b \to c$ transition as Cabibbo allowed because it is the dominant mode. The Cabibbo suppressed modes, in this terminology, are those amplitudes which are proportional to additional factors of sines.

where g_i are the weak couplings, we have

$$\langle f|T|i\rangle = g_1 M_1 e^{i\alpha_1} + g_2 M_2 e^{i\alpha_2}$$

$$\langle \bar{f}'|T|\bar{i}'\rangle = g_1^* M_1 e^{i\alpha_1} + g_2^* M_2 e^{i\alpha_2} \quad , \quad M_i = M_i^*$$

where the α_i are the strong interaction phases. Then the difference in the rates for $i \to f$ and $\bar{I} \to \bar{F}$ is given by

$$\Gamma - \bar{\Gamma} \sim \mathrm{Re}\left\{ (g_1^* g_2 - g_1 g_2^*) e^{i(\alpha_1 - \alpha_2)} \right\} \sim \mathrm{Im}(g_1^* g_2) \sin(\alpha_1 - \alpha_2)$$

We shall now proceed with estimation of the expected CP effects for the two-body decays of B_u and B_c mesons [$B_u^- = (b\bar{u})$, $B_u^+ = (\bar{b}u)$, etc.] taking the weak interactions at the tree level and assuming that the strong interactions are adequately described in terms of form factors. The transitions we have in mind are

$$b\,\bar{q}_1 \to (q_2\,\bar{q}_1)\,(d\,\bar{q}_2) \quad , \quad q_1 \neq q_2 \tag{3}$$

which, according to the conventional wisdom, take place via two separate mechanisms, the decay (W radiation) diagram (d) and the annihilation diagram (a) as depicted in Fig. 1. In reaction (3), q_j denotes a charge $\frac{2}{3}$ quark, u or c. The condition $q_1 \neq q_2$ is necessary in order to get different weak interaction phases in the two amplitudes.

The piece of the weak Hamiltonian responsible for reaction (3) is given by

$$H_W = \frac{G_F}{\sqrt{2}} \left\{ K_1 (\bar{d}_\beta c_\beta)(\bar{c}_\alpha b_\alpha) + K_2 (\bar{d}_\beta u_\beta)(\bar{u}_\alpha b_\alpha) \right\} + h.c.$$

$$K_1 = -s_1 c_2 [c_1 c_2 s_3 + s_2 c_3 e^{i\delta}]$$

$$K_2 = s_1 c_1 s_3 \tag{4}$$

where we have suppressed the V-A structure of the currents and used the common notation $s_i = \sin\theta_i$, $c_i = \cos\theta_i$; δ is the CP phase. Comparing with the diagrams in Fig. 1, one sees that for B_u decays the first term in H_W is responsible for the direct (radiation) amplitude and the second term for the annihilation process and vice versa, for the B_c decays. Note that the diagrams in Fig. 1 correspond to the vacuum insertion in the current-current product above.

3. - TWO-BODY DECAYS

We now explore CP effects in some two-body decays of B_u and B_c pseudoscalar mesons which can proceed through both mechanisms,

$$\begin{aligned} &\text{i)} \quad B_u^- \to D^- + D_o^* \\ &\text{ii)} \quad B_c^- \to \pi^- + \bar{D}_o \\ &\text{iii)} \quad B_c^- \to \pi^- + \bar{D}_o^* \end{aligned} \tag{5}$$

A fourth process

$$\text{iv)} \quad B_u^- \to D^- + D_o$$

proceeds only through the direct diagram and thus should show no CP violation if the vacuum insertion is a good approximation.

For the process i) the transition matrix element is of the form

$$\langle D^- D_o^* | H_W | B_u^- \rangle = \frac{G_F}{\sqrt{2}} (q_B \cdot \xi^*) \left\{ K_1 X^{(1)} + K_2 X^{(2)} \right\} \tag{6}$$

where q_B^μ is the four-momentum of B_u and ξ^μ is the polarization vector of D_o^*. The first term in (6) is due to the direct amplitude and the second one is the contribution of the annihilation mechanism. Furthermore,

$$(q_B \cdot \xi^*) X^{(1)} = \langle D^- | \bar{d} \gamma_\mu \gamma_5 c | 0 \rangle \langle D_o^* | \bar{c} \gamma^\mu \gamma_5 b | B_u^- \rangle$$

$$(q_B \cdot \xi^*) X^{(2)} = \langle D^- D_o^* | \bar{d} \gamma_\mu \gamma_5 u | 0 \rangle \langle 0 | \bar{u} \gamma^\mu \gamma_5 b | B_u^- \rangle \tag{7}$$

For the processes ii) and iii) we have similarly

$$\langle \pi^- \bar{D}_o | H_W | B_c^- \rangle = \frac{G_F}{\sqrt{2}} \left\{ K_2 \, Y^{(1)} + K_1 \, Y^{(2)} \right\} \tag{8}$$

$$\langle \pi^- \bar{D}_o^* | H_W | B_c^- \rangle = \frac{G_F}{\sqrt{2}} \left\{ K_2 \, Z^{(1)} + K_1 \, Z^{(2)} \right\} \tag{9}$$

where

$$Y^{(1)} = \langle \pi^- | \bar{d} \gamma_\mu \gamma_5 u | 0 \rangle \langle \bar{D}_o | \bar{u} \gamma^\mu b | B_c^- \rangle$$

$$Y^{(2)} = \langle \pi^- \bar{D}_o | \bar{d} \gamma_\mu c | 0 \rangle \langle 0 | \bar{c} \gamma^\mu \gamma_5 b | B_c^- \rangle \tag{10}$$

Finally $Z^{(i)}$ are obtained from $X^{(i)}$ by interchanging $u \leftrightarrow c$ and $\bar{u} \leftrightarrow \bar{c}$ in Eq. (7). Note that the matrix elements appearing in $X^{(1)}$, $Y^{(1)}$ and $Z^{(1)}$ are all directly measurable (at least in principle) in leptonic and semileptonic decays. In $X^{(2)}$, $Y^{(2)}$ and $Z^{(2)}$, however, the first factor is not directly accessible.

In order to give a theoretical estimate we need to know three types of matrix elements : axial current between pseudoscalar and vacuum, vector current between two pseudoscalars, and axial current between a pseudoscalar and a vector particle. The general form of these matrix elements is given by

$$\langle P(q) | A_\mu | 0 \rangle = f_P \, q_\mu$$

$$\langle P_2(q_2) | V_\mu | P_1(q_1) \rangle = F_+ (q_1 + q_2)_\mu + F_- (q_1 - q_2)_\mu$$

$$\langle V(q_2) | A_\mu | P(q_1) \rangle = \xi^{*\nu} \left[F_1 \, g_{\mu\nu} + F_2 \, q_{1\mu} q_{1\nu} + F_3 \, q_{2\mu} q_{1\nu} \right] \tag{11}$$

We use a mass breaking formula [10] for f_P with $f_\pi = 0.13$ GeV which leads to $f_D = 0.22$ GeV, $f_{B_u} = 0.36$ GeV, and $f_{B_c} = 0.41$ GeV. For the two-body matrix elements of the current, we make a naïve constituent quark model estimate with factorized flavour-spin and space-wave functions. The form factors $F((q_1-q_2)^2)$ then involve the characteristics of the quarks participating in the weak dynamics and the detailed momentum distribution provided by the meson wave functions. Since the meson wave functions are not known we follow other authors [11] and assume that the spatial wave functions for the initial and final mesons are the same. These inputs allow us to determine the normalization at $q^2 \equiv (q_1-q_2)^2 = 0$ and the relative q^2 dependence of the form factors. We find

$$F_-(q^2) = -\frac{M_1 - M_2}{M_1 + M_2} F_+(q^2) \quad , \quad F_+(0) = 1$$

$$F_2(q^2) = 0 \quad , \quad F_3(q^2) = \frac{-2}{(M_1+M_2)^2 - q^2} F_1(q^2) \quad , \quad F_1(0) = M_1 + M_2 \tag{12}$$

A difference between our results and earlier ones [10] is that we have kept the q^2 dependence of F_3/F_1 which follows from the model. Inserting (12) into the expressions for $X^{(i)}$, $Y^{(i)}$ and $Z^{(i)}$, $i = 1,2$, we find

i) $$X^{(1)} = f_D \; 2 M_{D_0^*} \frac{(M_{B_u} + M_{D_0^*})^2}{(M_{B_u} + M_{D_0^*})^2 - M_{D^-}^2} F_d^A (q^2 = M_{D^-}^2)$$

$$X^{(2)} = f_{B_u} \; 2 M_{D_0^*} \frac{(M_{D^-} + M_{D_0^*})^2}{(M_{D^-} + M_{D_0^*})^2 - M_{B_u}^2} F_a^A (q^2 = M_{B_u}^2)$$

ii) $$Y^{(1)} = f_\pi \frac{M_{B_c} - M_{D^0}}{M_{B_c} + M_{D_0}} \left[(M_{B_c} + M_{D^0})^2 - m_\pi^2 \right] F_d^V (q^2 = m_\pi^2)$$

$$Y^{(2)} = f_{B_c} \frac{m_\pi - M_{D^0}}{m_\pi + M_{D^0}} \left[(m_\pi + M_{D^0})^2 - M_{B_c}^2 \right] F_a^V (q^2 = M_{B_c}^2)$$

iii) $$Z^{(1)} = f_\pi \; 2 M_{D^{0*}} \frac{(M_{B_c} + M_{D^{0*}})^2}{(M_{B_c} + M_{D^{0*}})^2 - m_\pi^2} F_d^A (q^2 = m_\pi^2)$$

$$Z^{(2)} = f_{B_c} \; 2 M_{D^{0*}} \frac{(m_\pi + M_{D^{0*}})^2}{(m_\pi + M_{D^{0*}})^2 - M_{B_c}^2} F_a^A (q^2 = M_{B_c}^2) \tag{13}$$

Here F^V and F^A are the vector and axial form factors appearing in (12) but normalized to unity. The subscripts d and a denote the direct and the annihilation mechanisms. The general structure of the amplitudes in (13) is as expected on general grounds, in particular the appearance of the sum of the masses for $0^- \to 0^- + 1^-$ and the difference of the masses for $0^- \to 0^- + 0^-$. Note, however, that the dependence on M_B is vastly different for the annihilation contribution $Y^{(2)}$ than for $X^{(2)}$ or $Z^{(2)}$. This is simply due to the specific q^2 dependence of F_3/F_1 in (12) whereas $F_-/F_+ \sim$ constant is responsible for similar q^2 dependence of the coefficients in $Y^{(1)}$ and $Y^{(2)}$.

4. - RESULTS

Inserting the values of masses and f_P (P = π, D and B) we find

$$(2M_{D_0^*})^{-1} X^{(1)} = 0.24\ F_d^A(q^2 = M_{D^-}^2)\ ,\quad (2M_{D_0^*})^{-1} X^{(2)} = -0.53\ F_a^A(q^2 = M_{B_u}^2)$$

$$Y^{(1)} = 4.69\ F_d^V(q^2 = m_\pi^2)\ ,\qquad Y^{(2)} = 12.36\ F_a^V(q^2 = M_{B_c}^2)$$

$$(2M_{D_0^*})^{-1} Z^{(1)} = 0.13\ F_d^A(q^2 = m_\pi^2)\ ,\quad (2M_{D_0^*})^{-1} Z^{(2)} = -0.055\ F_a^A(q^2 = M_{B_c}^2) \tag{14}$$

in appropriate GeV units. We have taken $M_{B_u} = 5.02$ GeV and $M_{B_c} = 6.25$ GeV. Remembering that we must take the interference between the decay and annihilation contributions, we expect larger CP violating effects for processes i) and ii). Furthermore the time-like form factor $F_a(q^2)$ is probably larger at $q^2 = M_{B_u}^2$ than at $M_{B_c}^2$. Therefore, we give the explicit formulae for the first two cases :

i) $B_u^\mp \to D^\mp + \overset{(-)}{D}{}_0^*$

Here the relative difference between the conjugate decay modes, a manifestation of CP violation is given in our mechanism by the CP asymmetry parameter :

$$A^{CP} \equiv \frac{\Gamma^- - \Gamma^+}{\Gamma^- + \Gamma^+} = 2\, c_1 c_2 c_3\, s_2 s_3 s_\delta\, \mathrm{Im}\left[X^{(1)}X^{(2)*}\right] \Big\{ c_2^2 \left|c_1 c_2 s_3 + s_2 c_3 e^{i\delta}\right|^2 \left|X^{(1)}\right|^2$$

$$+ c_1^2 s_2^2 \left|X^{(2)}\right|^2 - 2 c_1 c_2 s_3 (c_1 c_2 s_3 + s_2 c_3 c_\delta)\, \mathrm{Re}\left[X^{(1)}X^{(2)*}\right] \Big\}^{-1},$$

$$s_\delta \equiv \sin\delta \;, \quad c_\delta \equiv \cos\delta \tag{15}$$

and $X^{(i)}$ are given in Eq. (14). The numerical value of the asymmetry depends on the parameters of the KM matrix as well as the form factors F_d^A and F_a^A at $q^2 = M_{B_u}^2$ which are all empirically unknown. We put the bottom charm form factor $F_d^A(q^2=M_{D^-}^2)$ equal to unity. The quantity F_a^A, which is an axial isovector charm-charm form factor is much more model-dependent. We shall estimate it in two rather orthogonal ways as follows.

Above 4 GeV, some knowledge [12)*)] of $\bar{D}D^*$ production in e^+e^-, dominated by ψ type vector isoscalar resonances, is available. If there were a broad strong axial resonance at those energies, its effects would still be noticeable at $q^2 = M_{B_u}^2$. Such a resonance with $M_A \simeq 4.2$ GeV and $\Gamma_A \simeq 300$ MeV would at $q^2 = (5.02\ \mathrm{GeV})^2$ give $F_a^A(q^2=M_{B_u}^2) = -2.3 + i(0.38)$. Note that in this case the annihilation contribution can in fact be larger than the direct one. Putting $c_1 = c_2 = c_3 = 1$, we find that the asymmetry is a function of the ratio $\lambda \equiv S_2/S_3$ and, of course, the phase δ. We have plotted the asymmetry in Fig. 2, as a function of λ, for several values of δ. The asymmetry is appreciable. For example, an asymmetry of about 30% would mean $\Gamma^- \sim 2\Gamma^+$, i.e., an appreciable difference in the partial branching ratios, $B_u^{\mp} \to D^{\mp} + \overset{(-)}{D}{}_o^{*}$.

If there are no resonances with the adequate quantum numbers we rely on perturbative QCD, neglect the mass effects and estimate the form factor following the procedure for the calculation [13)] of the pion form factor. In this case the absorptive part is much smaller (about 1%) and the asymmetry is given by

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*) We wish to thank Dr. R.H. Schindler for a very useful discussion on charm production.

$$A^{CP} \equiv \frac{\Gamma^- - \Gamma^+}{\Gamma^- + \Gamma^+} \simeq (2.2\,\%)\, R(\lambda, \delta) \tag{16}$$

$$R(\lambda, \delta) = \frac{2\lambda\, s_\delta}{1 + \lambda^2 + 2\lambda\, c_\delta} \quad , \qquad \lambda = \frac{s_2}{s_3} \tag{17}$$

$$R(\lambda^{-1}, \delta) = R(\lambda, \delta) \tag{18}$$

Here the 2.2% is simply $2.2\,\mathrm{Im}F_a^A(q^2=M_{B_u}^2)$. We have plotted the function $R(\lambda,\delta)$ in Fig. 3 for $\lambda \leq 1$. For $\lambda > 1$ one may simply use the relation (18). Note that in this case the annihilation contribution, being smaller than the direct one, has been neglected in the sum of the rates. For large values of δ the asymmetry is typically at the level of several per cent, again $\Gamma^- > \Gamma^+$.

The above results may be converted into the corresponding branching ratios using the CPT result $\Gamma(B_u^- \to \text{all}) = \Gamma(B_u^+ \to \text{all})$. An estimate for each individual partial rate can be obtained from the corresponding two-body Cabibbo allowed mode $B_u^- \to F^- + D_o^*$ which is expected [10] to be several per cent of the total rate. The ratio of interest is

$$\frac{\Gamma(B_u^- \to D^- + D_o^*)}{\Gamma(B_u^- \to F^- + D_o^*)} \sim s_1^2 \tag{19}$$

where s_1 is the Cabibbo angle. The exact expression for the ratio (19) is rather complicated and involves the parameters of the KM matrix, etc.

ii) $B_c^\mp \to \pi^\mp + \overset{(-)}{D}{}^o$

The CP asymmetry parameter in this case is obtainable from the invariant amplitudes $Y^{(1)}$ and $Y^{(2)}$ in Eq. (14). The annihilation πD form factor now enters at $M_{B_c}^2$ which is far above the expected resonances with charm quantum numbers. Therefore we neglect the annihilation contribution in the sum of the rates and find

$$A^{CP} \equiv \frac{\Gamma^- - \Gamma^+}{\Gamma^- + \Gamma^+} \simeq -2\lambda\, s_\delta\, [2.6\ \mathrm{Im}\, F_a^V(q^2 = M_{B_c}^2)] \;, \quad \lambda = \frac{s_2}{s_3} \tag{20}$$

The branching ratio as compared to the corresponding two-body mode $B_c^- \to \pi^- + \eta_c$ is given by

$$\frac{\Gamma(B_c^- \to \pi^- + \bar{D}_o)}{\Gamma(B_c^- \to \pi^- + \eta_c)} \simeq \frac{s_1^2}{1 + \lambda^2 + 2\lambda c_\delta}$$

where the $D^o - \eta_c$ mass difference has been neglected.

We note that the CP asymmetry depends on quite different combinations of the KM parameters in the two processes i) and ii). Here the effect (20) is linear in λ and in $\sin\delta$, in contrast with those shown in Figs. 2 and 3 for the case i). The order of magnitude of the effect is proportional to the absorptive vector form factor F_a^V at $q^2 = M_{B_c}^2$. It is plausible that since one is far above the relevant resonance region with charm quantum number, the asymptotic form factor predicted [13] by QCD should give a reasonable estimate. We apply it and find $\mathrm{Im}\, F_a^V(q^2 = M_{B_c}^2) \simeq 5\times10^{-3}$. Insertion into Eq. (20) yields

$$\mathcal{A}^{CP} \equiv \frac{\Gamma^- - \Gamma^+}{\Gamma^- + \Gamma^+} \simeq 3 \left(\frac{s_2}{s_3}\right) s_\delta \ \% \tag{21}$$

5. - CONCLUSIONS

To summarize, we have shown that for charged bottom meson decays, CP violation is induced through a mechanism of interference between two different tree-level charged current amplitudes. They correspond to the direct (decay) b quark diagram and the annihilation diagram. The vertices associated with them provide a relative observable phase within the scheme of the standard electroweak theory with three families of quarks. It is thus predicted that direct CP violation in the transition matrix of bottom particles is present. The Cabibbo disfavoured modes enhance the CP violating asymmetry, which can be maximal at the level of the couplings. The asymmetry in the rates of conjugate decay processes has been considered and applied to two-body decays of B_u and B_c mesons. A realistic calculation of the effects depends on dynamics of non-leptonic interactions which are not well known. With a model-dependent calculation of the weak hadronic vertices, inserting the vacuum into the current-current product, we find that the process $B_u^- \to D^- + D_o^*$ is promising for the effects discussed here. The results are given in Figs. 2 and 3 [Eq. (16)], for two orthogonal views of the annihilation axial form factor, as a function of the KM angles and phase parameters. The effect ranges from a few per cent to very big values 50%. For $B_c^- \to \pi^- + \bar{D}_o$ the annihilation contribution is expected to be small, so the CP effect (20) is probably at most several per cent. These estimates suggest that there could be an important direct CP violation in B decays. If observed, it will provide an essential clue to the understanding of the origin of CP violaton. The extension of the mechanism discussed in this paper to semi-inclusive decays of the B mesons will be treated separately.

ACKNOWLEDGEMENTS

We wish to thank the CERN Theoretical Studies Division for the hospitality extended to us. We have enjoyed discussing this work with A. De Rujula, C. Sachrajda and F.J. Yndurain.

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FIGURE CAPTIONS

Figure 1 Direct decay (d) and annihilation (a) diagrams for $(b\bar{q}_1) \to (d\bar{q}_2)(q_2\bar{q}_1)$.

Figure 2 CP asymmetry between the conjugate decay rates for $B_u^- \to D^- + D_o^*$, as function of $\lambda \equiv s_2/s_3$, for different values of the phase δ. An annihilation axial form factor dominated by a strong resonance has been assumed (see the text).

Figure 3 The function $R(\lambda \equiv s_2/s_3), \delta)$ defined in Eq. (17), relevant for the CP asymmetry of $B_u^- \to D^- + D_o^*$ in the case of a small annihilation form factor.

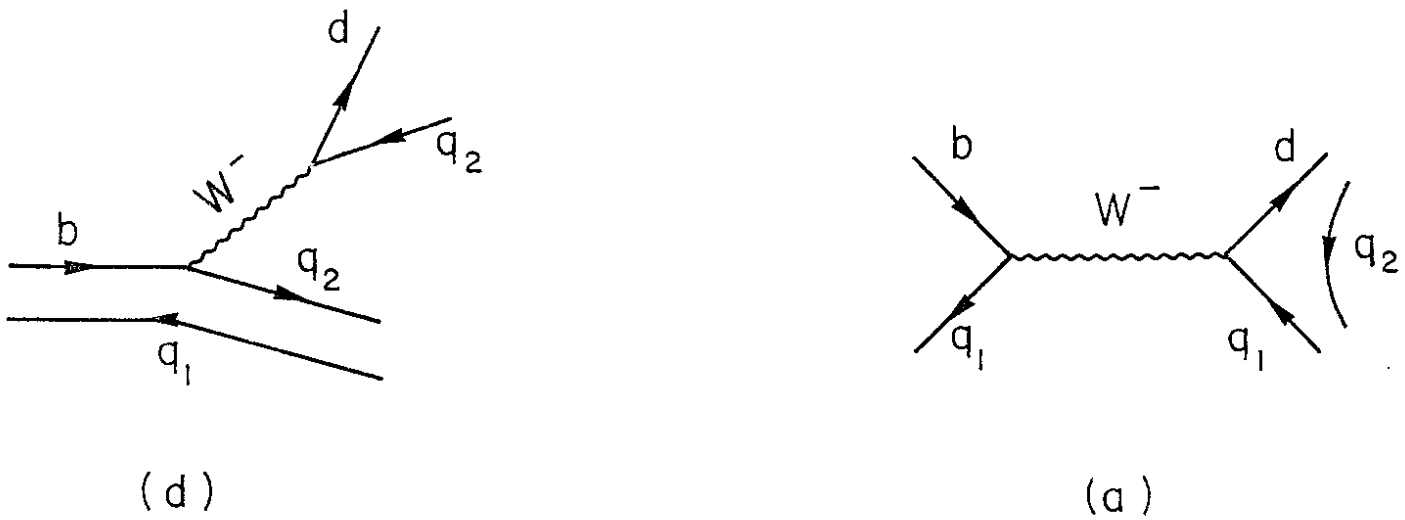


Fig. 1

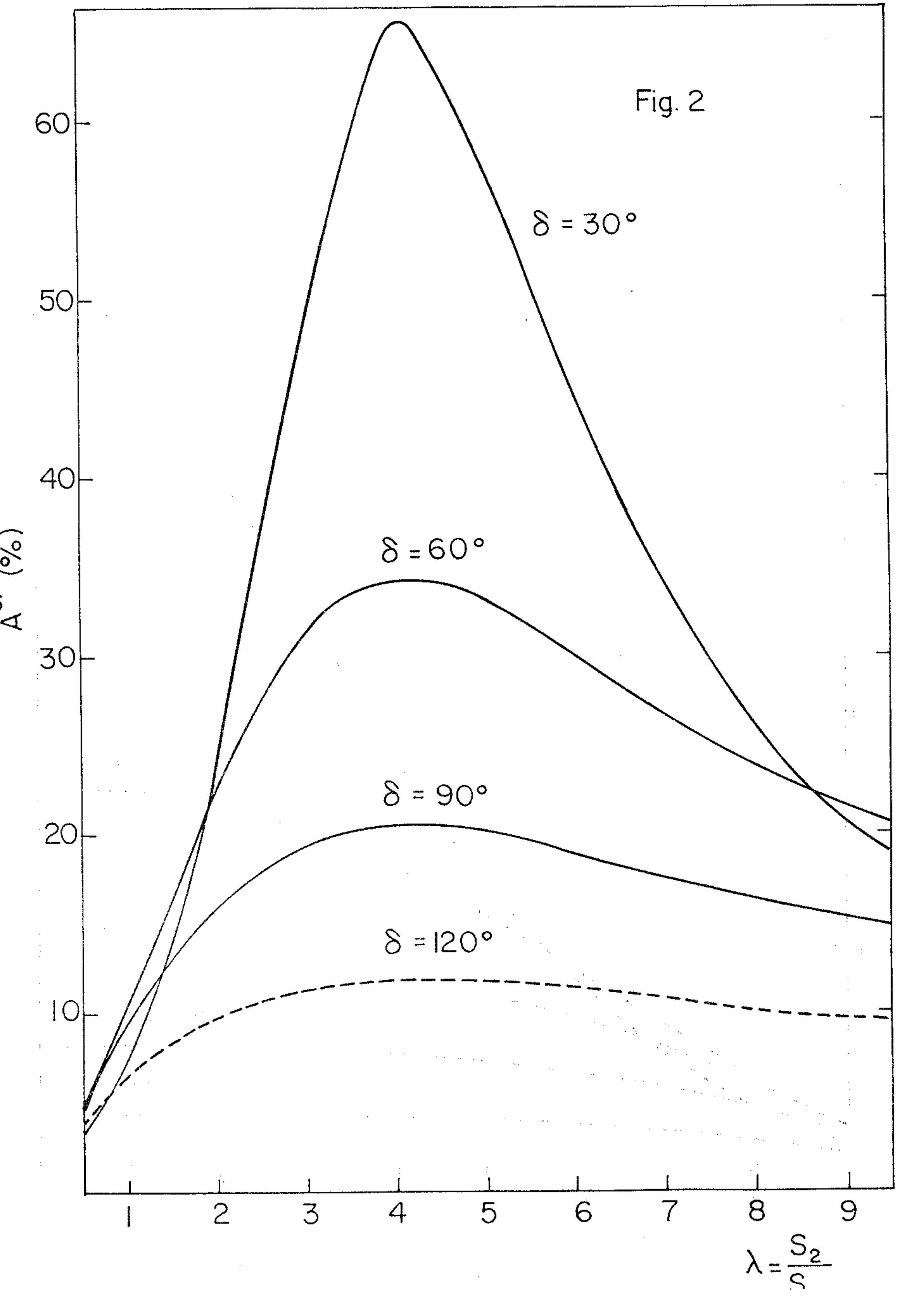
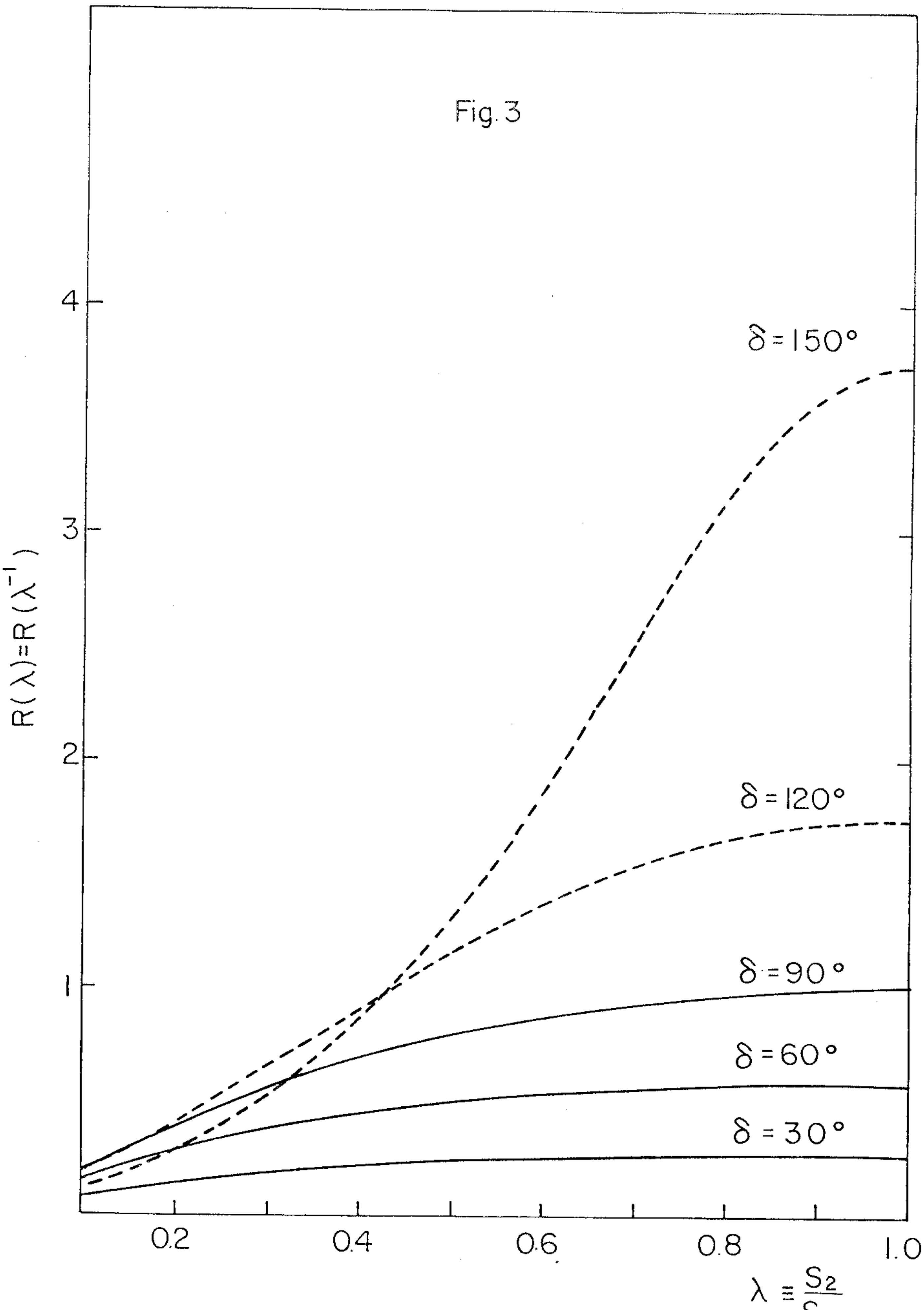


Fig. 2
$\delta = 30°$
$\delta = 60°$
$\delta = 90°$
$\delta = 120°$
A^{CP} (%)
60
50
40
30
20
10
1
2
3
4
5
6
7
8
9
$\lambda = \frac{S_2}{S}$

Fig. 3



4
3
2
1
$R(\lambda)=R(\lambda^{-1})$
$\delta = 150°$
$\delta = 120°$
$\delta = 90°$
$\delta = 60°$
$\delta = 30°$
0.2
0.4
0.6
0.8
1.0
$\lambda \equiv \frac{S_2}{S_-}$

18.