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OF THE ARGUS MEASUREMENTS

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PHYSICS OF TOP AND CP VIOLATION IN B DECAYS
IN THE LIGHT OF THE ARGUS MEASUREMENTS

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ABSTRACT

Implications of the significant $B_d^0-\bar{B}_d^0$ mixing, observed by the ARGUS collaboration, on our present understanding of nature are discussed. Our emphasis is on the consequences of the standard model with a heavy top quark. We discuss also the future prospects of the observation of CP violation in B decays.

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1. INTRODUCTION

According to the Greek mythology the vigilant giant ARGUS has been put on guard of the beautiful Io. To free Io from the ARGUS surveillance he has been lulled by the sweet talks. As if to confirm the mythological origin of the acronym of its name, just to the 10th anniversary of the discovery of beauty the ARGUS collaboration has made the sensational observation of the large B^0 \bar{B}_d^0 mixing [1] and of the charmless B decays [2,3].

These fundamental results, assuming they are confirmed, keep the key place in the ranks of the recent experiments [4] (new UAI limit for m_t , nonvanishing value for $(\epsilon/\epsilon)_K$, etc) providing the vigilant surveillance of the phenomenological position of the standard model (SM). As is well known nowadays the SM has achieved the overwhelming phenomenological successes describing the enormous wealth of the experimental data (the so called festival of the SM). Our purpose here is to demonstrate how the ARGUS results constrain the degrees of freedom of the SM and to enumerate the most important implications, namely, restructuring of our understanding of the physics of top, prospects of observing the CP violation effects in B decays and one of the consequences for the Higgs sector.

2. THE FORTUNE OF THE TOP QUARK [5]

Of the six personages in the minimal scenario of the SM five have been observed so far. We can feel quite confident that the sixth flavour - top - exists as well. Within the SM the ARGUS finding requires a heavy top quark $m_t \approx 100 - 150 \text{ GeV}^2$ (more conservatively $m_t \gtrsim 60 \text{ GeV}$ accounting for the existing uncertainties in f_B , x_{B_d} and R, see below). A comprehensive analysis of electroweak phenomena leads to the restrictions [4,8] $44 \text{ GeV} \lesssim m_t \lesssim 190 - 210 \text{ GeV}$. The properties of the ultraheavy top should be in the marked contrast to the case of its predecessors - c and b; in some sense they are expected to be closer to those

2) The literature in this area has become quite vast, see e.g. review [6]. We shall follow here Ref[7]

of heavy leptons (see Refs[9,10,5]). In the first place this difference is connected with the large t-quark width [9]. At $m_t > M_W$ t-quark decays semiweakly $t \rightarrow b + W^+$. In the SM the width $\Gamma_t \equiv (t \rightarrow b + W^+)$ quickly reaches its asymptotic form ($|V_{tb}| \approx 1$)

$$\Gamma_t \approx 175 \text{ MeV} (m_t/M_W)^3 \quad (1)$$

(for the exact calculations of the t quark width see Ref[9]). At $m_t > 100 \text{ GeV}$ the t lifetime $\tau_t = \Gamma_t^{-1}$ becomes shorter than the typical strong interaction time scale $t_{\text{form}} \sim R \sim \Lambda^{-1}$ (the time needed to build up (tq), (tq) hadrons by picking up light quarks),

$$\tau_t \sim 1 \text{ fm} (M_W/m_t)^3 < t_{\text{form}} \sim \Lambda^{-1} \sim 1 \text{ fm}.$$

This prevents the bound states from being forming. For such a quark all the bremsstrahlung processes prove to be under the jurisdiction of perturbative QCD. One can say that this quark in many aspects behaves as if it were a free coloured object (see Ref[10]). In particular, t quark is not depolarized through the fragmentation and the gluon bremsstrahlung effects.

Let us come to the discussion of the threshold behaviour expected for heavy top production in e^+e^- collisions. As it is well known [11,12] for large enough quark masses the influence of the nonperturbative effects on the properties of the bound states with lower quantum numbers is small, and the quark dynamics is governed only by the electroweak and perturbative QCD interactions. In the threshold region the interaction is determined by the two-particle Coulomb-like potential. Owing to the multiple soft gluon exchange the quarks are bound at the distances $Z \sim k^{-1}$ ($\alpha_s m_t$)⁻¹ and (at least for the lowest levels) the standard set of the Coulomb-like bound states is formed below the continuum threshold (see Refs[11-13]).

For the lowest vector state - toponium $T(1^3S_1(t\bar{t}))$ the characteristic size of the $t\bar{t}$ bound system is $Z_f \sim k_1^{-1}$

$$k_1 = 2/3 m_t \alpha_s(k_1),$$

the level spacing being

$$2m_t - M_{10} = \Delta E_1 = 4/9 m_t \alpha_s^2(k_1). \quad (2)$$

The account for the Coulomb effects and for the large \sqrt{t} modifies drastically the threshold behaviour of $\sigma_{t\bar{t}}$. As the result one should replace the standard threshold factor $Z^{(0)} = 3/2 \beta_t$ ($\beta_t = \sqrt{1-4m_t^2/s}$ is the centre-of-mass velocity of the top quark), corresponding to the vector contribution into $\sigma_{t\bar{t}}$, by

$$Z_v^{(0)} \rightarrow Z_v = 3/2 \left(\frac{4\pi}{m_t^2} \text{Im} G_{E+1} \Gamma_t(0,0) \right). \quad (3)$$

Here $G_E(\vec{z}, \vec{z})$ is the Green function of the $t\bar{t}$ system in the colour singlet state, $E = \sqrt{s} - 2m_t$ - nonrelativistic energy of quarks. The axial vector contribution is evidently small throughout the threshold region: $\mathcal{Z}_A \sim \pi \alpha_s \beta_t^2$. Note that the account for \sqrt{t} is performed by replacing E by $E+i\Gamma_t$. One may be easily convinced in this using Coulomb gauge. The Coulomb effects lead to a sharp increase of $\sigma_{t\bar{t}}^{\text{non}}(e^+e^- \rightarrow t\bar{t})$ in the continuum region. Thus, e.g., in the narrow width approximation

$$Z_v^{(0)} \rightarrow Z_v = \frac{3}{2} \beta_t \frac{Z}{1 - e^{-Z}}, \quad Z = \frac{4}{3} \frac{\pi \alpha_s}{\beta_t},$$

(see Refs[5,13]) and at $\beta_t \rightarrow 0$ the threshold factor transforms to the finite quantity $Z_v = 2\pi \alpha_s$. At fixed colour coupling α_s the explicit form for $\text{Im} G_{E+1}^{(0)}$ is given by formula

$$\text{Im} G_{E+1}^{(0)}(0,0) = \frac{m_t^2}{4\pi} \left[\frac{k_+}{m_t} + \frac{2\bar{K}_1}{\beta_t} \arctan \left(\frac{k_+}{k_-} \right) \right] + \sum_{n=1}^{\infty} \frac{2\bar{K}_1^n}{m_t^n} \frac{\Gamma_t \bar{K}_1^n + k_+ (n^2 \sqrt{E^2 + \Gamma_t^2} + \bar{K}_1^2/m_t)}{(E + \bar{K}_1^2/m_t n^2)^2 + \Gamma_t^2}; \quad \bar{K}_1 = \frac{2}{3} m_t \alpha_s,$$

$K_+ = \sqrt{m_t^2 + \sqrt{E^2 + \Gamma_t^2}}$. The first term in the square brackets in Eq.(4) corresponds to the Born approximation, modified by the effects of \sqrt{t} , the second term corresponds to the one-loop correction and the third one - to the sum over the bound S-wave states, that acquire the widths $\Gamma = 2\Gamma_t$. We shall use below the eq.(4) to illustrate the

evolution of the $e^+e^- \rightarrow t\bar{t}$ reaction in the threshold region with the increasing of the top mass m_t .

With the account for the QED effects³⁾ the expression

$$\text{for } R_{t\bar{t}} = \sigma_{t\bar{t}} / \sigma_{pt}$$

$$(\sigma_{pt} = 4\pi\alpha_s^2/3s) \text{ takes the form}$$

$$R_{t\bar{t}} = \frac{9}{2} \cdot \frac{4\pi/m_t^2}{\alpha} \left\{ Q_t^2 + \frac{v_t^2}{\alpha(1-M_W^2/4m_t^2)} \right\} \frac{(1+3/4\beta)}{(1-P(4m_t^2))} \cdot \int_0^1 dx \beta^{1-x} \text{Im} G_{E-m_t x + i/\Gamma_t}(0,0), \quad (5)$$

where $Q_t = 2/3$, $v_t = 1 - 8/3 \sin^2 \theta_W$ is the t-quark vector constant; $\alpha = (16 \sin^2 \theta_W \cos^2 \theta_W)^{-1}$; ($\alpha = 8$ at $\sin^2 \theta_W = 0.23$); $\beta = \frac{4s}{W}$. ($\ln 2m_t/m_e \rightarrow 1/2$), $P(4m_t^2)$ is the real part of the photon vacuum polarization operator.

An analysis of the threshold behaviour of the $e^+e^- \rightarrow t\bar{t}$ reaction leads to the following conclusions (see for details Ref[5])

1. At $m_t < 100$ GeV the standard set of the Coulomb-like bound states is formed. The ground state width $\Gamma_T = 2\Gamma_t \approx 170$ MeV at $m_t = 100$ GeV, so that the toponium and the continuum are rather well separated: $\Delta E_1 \approx 1$ GeV at $\alpha_s = 0.150$, see Eq(2). The cross section in the T-resonance maximum with the account for the QCD and electroweak effects is given by

$$R_T = \frac{\sigma(e^+e^- \rightarrow T \rightarrow t\bar{t})}{\sigma_{pt}} \approx \frac{2\alpha_s^3(k_1/m_t)}{\Gamma_T} = \frac{9}{2} \alpha_s(k_1) \frac{\Delta E_1}{\Gamma_T},$$

where Γ_T is the total width of toponium $\Gamma_T \approx 2\Gamma_t$ for $m_t > M_W$.

3) Note that the QED radiative corrections connected with photon bremsstrahlung from the initial electrons are very important here: the height of the resonance curve decreases by a factor of two and the so-called radiative tail arises, see, e.g. Ref[14]

2. With the increasing of m_t the heights of the resonance curves decrease, their widths increase, and the continuum is shifted to the left rather rapidly owing both to the finite width and Coulomb effects. At $m_t > 150$ GeV the decay lifetime of toponium τ_T becomes shorter than the Bohr orbit period, τ_B , of the toponium.

$$\tau_T = \frac{1}{\Gamma_T} \approx \frac{1}{2\Gamma_t} < \tau_B \approx \frac{1}{\Delta E_t} \sim 0.13 \text{ fm},$$

(see Ref[9]). Thus, the toponium bound state actually can no longer be formed.

3. With the further increase of m_t the $t\bar{t}$ levels completely overlap, and at $m_t \approx 200$ GeV ($\Delta E_1 = 1.5$ GeV at $\alpha_s = 0.13$, $\Gamma_T \approx 5$ GeV) all resonance structures are practically washed out and buried in the continuum. All the properties of the process $e^+e^- \rightarrow t\bar{t}$ would follow literally the predictions of the free quark model, modified only by perturbative QCD corrections.

Analogously for the massive quarks in the case of extra generations the quarkonium states can form only if the semiweak decay is suppressed in some way.

In hadron collisions the threshold behaviour of t, \bar{t} production is dominated by the S-wave colourless states of the quark pair, and quarkonium production is only a very small fraction of the $t\bar{t}$ cross section. The Coulomb effects may be accounted for analogously to Eq(4).

Note that the analysis in Ref[5] in many aspects was performed only on the illustrative level.

For more comprehensive quantitative analysis it is necessary to account for consistently the running α_s or to find $G_{E+i\Gamma_t}(0,0)$ not for the Coulomb potential, but for the potential $\mathcal{V}_t(r)/2$. In addition one should account for the dependence of the top quark width on its virtuality, this dependence proves to be essential in the region $|E| \gtrsim m_t - m_b - M_W$ that is near the threshold at $m_t < 100$ GeV.

At least for $(E^2 + \Gamma_t^2)^2 < \alpha_s^2 < \frac{\alpha_s^2}{\pi} < \alpha_s^2$ (which exists at $m_t < 100$ GeV the correct account for the nonperturbative effects is needed.

The uncertainties of calculations decrease with m_t increasing.

Thus the threshold behaviour of $t\bar{t}$ production is completely determined by a calculable function of m_t , Γ_t and α_s , whose shape changes with the variation of m_t . An experimental study of the threshold behaviour would permit one to measure accurately m_t and Γ_t (and therefore to test the SM prediction $|V_{tb}| \approx 1$ and, in principle, would enable us to determine α_s under the most favourable theoretical conditions.

Note also that it is important from the point of view of testing SM to compare the top data with the results of the precise measurements of M_Z . The point is that in the framework of SM the allowed regions of m_t and M_Z are correlated and the deviations may indicate in favour of the new physics.

A serious background to the processes $e^+e^- \rightarrow t\bar{t}$ at $m_t > M_W$ arises from the W^+W^- pair production, $e^+e^- \rightarrow W^+W^-$. The corresponding cross section is relatively large and approximately flat in the energy range $\sqrt{s} \approx 200-300$ GeV

$$R_{W^+W^-} = \sigma(e^+e^- \rightarrow W^+W^-) / \sigma_{pt} \sim 10.$$

This background can be substantially reduced using longitudinally polarized initial particles. The point is that the main role in the process $e^+e^- \rightarrow W^+W^-$ is played by the neutrino exchange graph, which contributes only to $\mathcal{O}(e_2^+e_2^-)$ due to the V-A structure of the charged current interactions.

Finally, considering all the features of the ultraheavy top quark, it seems hard to imagine that nature would deny us such a nice (but scanty) laboratory.

3. CP VIOLATION IN BEAUTY DECAYS⁴⁾

As is well known the theorists remain unhappy in the face of such a success of the SM, because it is not a fundamental theory and is thought to be only an incomplete reflection of nature.

To lull the ARGUS vigilance and, thus, to free the theory it is not enough to listen to the sweet theoretical talks on the possible extensions of the SM (extended gauge models, SUSY, compositeness, etc). Further experimental efforts to search for new physics are badly needed.

The examination of CP violating phenomena seems to be one of the major milestones on the so called "low road to new physics" (i.e. the indirect search of the manifestations of new forces), see e.g. Ref[16]. We feel strongly that CP violation has to be found outside the decays of kaons before light can be shed on the source of this phenomena. As is well known beauty particles carry nice promise to exhibit large CP violating effects in their decays.

In the SM CP violation comes from the one irreducible phase in the Kobayashi-Maskawa (KM) [18] quark mixing matrix.

A simple way exists to represent geometrically the CP-odd phases of the KM matrix elements V_{ij} without appealing to any particular parametrization of the mixing matrix.⁵⁾ Indeed, the unitarity of the KM matrix yields, among others, the relation

$$V_{ub}^* V_{ud} + V_{cb}^* V_{cd} + V_{tb}^* V_{td} = 0, \quad (6)$$

4) We shall follow here paper [15] dealing mainly with e^+e^- collisions (see also Refs[7,16]). There is a number of excellent reviews, discussing the prospects for B physics at hadronic machines (see e.g. Ref[17])

5) To our knowledge this geometrical approach was first pointed out by Bjorken, L.-L. Chau and Jarlskog and became nowadays a folklore. The present authors independently used it in their lectures at the LNPI and ITEP schools, see also Ref[16]

which is evidently invariant under changes in the phase convention adopted for the quark fields.

This equation represents triangle relation in the complex plane (as shown in Fig. 1a), and is quite accessible to geometrical intuition. Different parametrizations of the KM matrix correspond just to some rotations of this triangle as a whole.

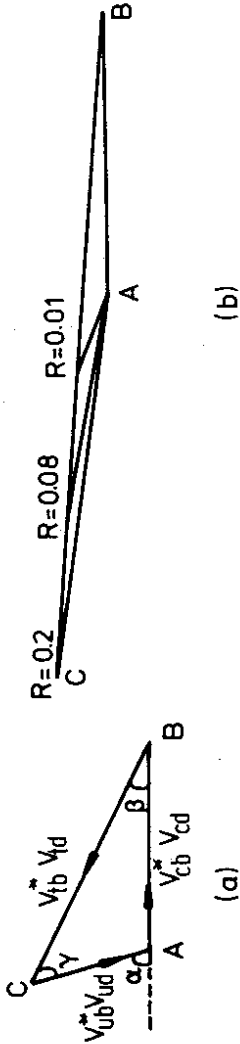


Fig.1

Fig. 1 The unitarity triangle in KM ansatz with three families.

It is important that in many cases the features of B physics, and in particular, the CP-odd observables, have the explicit geometrical interpretation in terms of this "unitarity triangle" (UT). Thus, all the main CP-odd phases for B decays are just the angles of UT. Note, that the side AB is more or less known, because $|V_{cd}| = |V_{us}| = \sin c \approx 0.221$, and $|V_{cb}|$ can be extracted from the beauty lifetime

$$|V_{cb}|^2 \sim \frac{1}{\tau_B} \cdot \frac{1}{1+R} \quad (7)$$

here R is

$$R \approx \frac{\Gamma_{se}(b \rightarrow u)}{\Gamma_{se}(b \rightarrow c)} \approx 2.1 \left| \frac{V_{ub}}{V_{cb}} \right|^2$$

(The numerical factor accounts here for the phase space difference in the $b \rightarrow u$ and $b \rightarrow c$ decays).

The length of the side AC defines the amplitudes of the KM-suppressed decays of beauty in the $b \rightarrow u$ transition. If one

normalizes the overall scale as if to have the known side AB to be of unit size, then the length of the AC side would be

$$\left| \frac{AC}{AB} \right| \approx \left| \frac{V_{ub}}{V_{cb}} \right| \cdot \text{tg } \theta_c \approx \sqrt{10R} \quad (8)$$

The third side, BC, is connected with the $B_d - \bar{B}_d$ mixing. For the known values of the m_c and of the hadronic matrix elements the length of the BC side can be expressed through the mixing parameter $X_{B_d} = \left(\frac{\Delta m}{\Gamma} \right)_{B_d}$. For the "normalized" UT the length is obviously equal to $|1 + \sqrt{10R} e^{i\omega}|$ (see Fig. 1a).

The UT is actually unique. The triangles, corresponding to other orthogonality conditions, are equal to it (up to small corrections $\sim \sin^2 \theta_c$, $|V_{cb}|^2$) or are quite squashed ones. For example, the orthogonality condition for the first and the third rows of the KM matrix leads to

$$V_{ub}^* V_{ts} + V_{us}^* V_{td} + V_{ud}^* V_{td} = 0 \quad (9)$$

It is easy to see that up to terms $\sim \sin^2 \theta_c$, $|V_{bc}|^2$ the triangle, corresponding to Eq(9), is equal to the original one and they can be rotated from one to another. Thus, for known value of $X_{B_d} \approx 0.73$ the measurement of any two quantities of m_c , R or angles of the UT would completely define it, and the further accurate measurements enable us to check the SM prediction Eq(6), i.e. the closeness of the triangle.

The mixing parameter X_{B_d} is given by

$$X_{B_d} \approx 0.10 \left(\frac{m_t}{80 \text{ GeV}} \right)^2 \cdot \frac{\xi (m_t^2/M_W^2)}{3/4} \cdot \frac{|1 + \sqrt{10R} e^{-i\omega}|^2}{1+R}$$

where $\xi (m_t^2/M_W^2)$ is the known W propagator factor ($\xi(1) = 3/4$), and we use the value $f_B = 100 \text{ MeV}$.

Up to now rather scanty information on V_{ub} is available: $0.01 \leq R \leq 0.08 - 0.5$ (see e.g. [3]). The upper bound $R < 0.08$ comes from the leptonic spectra in semileptonic B decays. But this analysis

is highly model dependent. The charm counting in B decays is reliable, but it leads now only to a rather weak bound $R < 0.2 - 0.5$. At the same time the ARGUS finding of the charmless decays $B \rightarrow \bar{p}\bar{p}\pi^0$ [2] indicates that the ratio R is not too small, $R \gtrsim 0.05 - 0.1$ for average values; some model estimates prefer even $R > 0.2$. The phases between the different KM matrix elements represent the only source of the CP violation in the SM, therefore they can be found from the CP violation parameter for K_L^0 decay, ξ_K .

Let us formulate the results of the combined analysis for the UT, accounting for the ARGUS results (see also Ref[7]).

1. All the angles α , β and γ in the UT lie in the first quadrant and are quite small. It means that the triangle has, indeed, the topology shown in Fig. 1a, but it is almost degenerate. Therefore the CP violation is far from being maximal one.

2. The angle β appears to be remarkably stable with respect to variation of the poorly known quantity R (the side AC). Numerically it is about $\beta \approx 0.075 \approx 4.4^\circ$ (for $f_B = 100$ MeV, $X_{B_d} = 0.73$, $B_K = 1$; see below). The real shape of the UT is shown in Fig. 1b for $R = 0.2$; 0.08 and 0.01. Note that at $R \approx 0.1$ the UT is isosceles, hence $\gamma = \beta$ and $\alpha = 2\beta$.

Of course, one is not able now to predict quite accurately the CP odd angles of the triangle, despite the stability of β mentioned above. Actually the calculation includes the hadronic matrix elements for both B and K mesons, which are estimated theoretically; the angles depend also on the exact values of X_{B_d} and \mathcal{T}_B . The main uncertainty is due to the value of f_B (we assume $B_B = 1$; in fact the relevant combination is $B_B f_B^2$), B_K and $|V_{cb}|^2$. The angle β and, consequently, α and γ (for $R > 0.01$ when α and γ are still small) appear to depend as

$$\beta \sim \frac{f_B}{X_{B_d}} \cdot B_K^{-1} \tau_B^2 \quad (11)$$

at fixed value of R. Eq(11) enables one to see easily the variation of the CP odd effects in B mesons when different parameters are changed.

Let us turn to the expected CP asymmetries in B mesons. First of all we shall discuss the difference in same sign dileptons, arising from the production of $B^0 \bar{B}^0$ pairs

$$A = \frac{N(e^+e^+) - N(e^-e^-)}{N(e^+e^+) + N(e^-e^-)} \quad (12)$$

This asymmetry is purely a "superweak" effect, i.e. it is completely determined by the CP violation in the $B^0 - \bar{B}^0$ mixing ($A = \frac{4\text{Re}\xi_B}{1 + |\xi_B|^2}$). The latter appears when the absorptive, $\sqrt{S}\xi_B$, and the dispersive, $M_{B\bar{B}}$, parts of the total transition amplitude $B^0 \rightarrow \bar{B}^0$ have the different phases. The direct calculation shows that, if \sqrt{S} is described by the standard quark diagrams, the asymmetry A is proportional to the area of the UT for both B_d and B_s mesons:

$$A_{B_d} \approx -10 \left(\frac{f_B}{100 \text{ MeV}} \right)^2 \cdot \frac{S}{X_{B_d}} \approx - \frac{1.2 \cdot 10^{-3} \left(\frac{f_B}{100 \text{ MeV}} \right)^2 S}{1 + R} \quad (13)$$

$$A_{B_s} \approx 10 \left(\frac{f_{B_s}}{100 \text{ MeV}} \right)^2 \cdot \frac{S}{X_{B_s}} \approx \frac{1.2 \cdot 10^{-3} \left(\frac{f_{B_s}}{100 \text{ MeV}} \right)^2 S}{1 + R}$$

where S is the area of the triangle while \tilde{S} is the area of the normalized UT with $AB = 1$. Therefore one has in the SM

$$\frac{A_{B_s}}{A_{B_d}} \approx - \left(\frac{X_{B_d}}{X_{B_s}} \right) \left(\frac{f_{B_s}}{f_{B_d}} \right)^2 \quad (14)$$

It follows from Eqs(13) that if $S, \tilde{S} \sim X_{B_d}^{-1}$, the "superweak CP violating effects should decrease with increasing of X_{B_d} as $A \sim X_{B_d}^{-2}$. This explains why the ARGUS finding $X_{B_d} \approx 0.73$ leads to the considerable decrease of the expected CP asymmetry. Numerically the predictions for A which follow from Eqs(13) are

$$A_{B_d} \approx -2 \cdot 10^{-4}, \quad A_{B_s} \approx 4 \cdot 10^{-5} \quad (15)$$

for $R \approx 0.1$. Note that A_{B_d} decreases slightly with R decreasing. So in the SM one needs at least 10¹¹ $B\bar{B}$ events to observe such asymmetry. New physics might well enhance the leptonic

asymmetries, and it seems to be quite reasonable to look at them in future experiments.

The study of the "milliweak" CP nonconservation, which generally does not require the CP violation in mixing, looks much more promising. Here the certain decays of the neutral B mesons, where the CP odd effects result from the interference of the decay channels $B^0 \rightarrow f$ and $B^0 \rightarrow \bar{B}^0 \rightarrow f$, are expected to be the most convenient owing to the largest predicted effects and to the unambiguity of the theoretical interpretation. The latter is particularly clear for the decays to the states with definite CP parity χ_f . For such decays the CP odd asymmetries depend on one CP violating phase due to the complexities of the relevant KM matrix elements. For example,

$$\Delta = \frac{\Gamma(B^0 \rightarrow f) - \Gamma(\bar{B}^0 \rightarrow \bar{f})}{\Gamma(B^0 \rightarrow f) + \Gamma(\bar{B}^0 \rightarrow \bar{f})} = -\chi_f \sin 2\phi \frac{x}{1+x^2} \quad (16)$$

In the SM there are three main types of B decays where one can predict the sizeable CP asymmetry:

- (i) The allowed decay of B_d via the transition $b \rightarrow c\bar{c}s$, e.g. $B_d(\bar{B}_d) \rightarrow J/\psi K_S$ (Cabibbo-suppressed channel $b \rightarrow c\bar{c}d$, e.g. $B_d(\bar{B}_d) \rightarrow D^+ D^-$, $J/\psi \pi^0$).
- (ii) The KM-suppressed B_d decays via the transition $b \rightarrow u\bar{u}d$ like $B_d(\bar{B}_d) \rightarrow \pi^+ \pi^-$.
- (iii) The KM-suppressed decays of $B_s(\bar{B}_s)$ via the transition $b \rightarrow u\bar{u}d$, say $B_s \pi^0 K_S$.

The relevant phases ϕ are just the angles of the UT:

- (i): $B_d, b \rightarrow c\bar{c}s (b \rightarrow c\bar{c}d)$ $\phi = -\beta$,
- (ii): $B_d, b \rightarrow u\bar{u}d$ $\phi = \gamma$,
- (iii): $B_s, b \rightarrow u\bar{u}d$ $\phi = \alpha$,

see Fig. 1a. For the allowed decays of B_s via $b \rightarrow c\bar{c}s$ (and for $b \rightarrow c\bar{c}d$) the corresponding phase is small, $\phi \sim \sin^2 \theta_c \sqrt{10R} \sin \alpha$.

Somewhat different situation occurs for the decays like $B_d^0 \rightarrow D^+ \pi^-$ or $D^+ \pi^-$, where the transition $b \rightarrow c\bar{u}d$ and $\bar{b} \rightarrow \bar{u}c\bar{d}$ interfere. Here the CP odd phase difference 2ϕ equals to $(\gamma - \beta)$, in other words ϕ is the angle between the external bisector \vec{s} of the angle CAB and the side BC, see Fig. 2. It is evident that at $R \approx 0.1$ when the UT is isosceles, the CP odd effect in such processes vanishes; at other realistic values of R this phase remains suppressed being of opposite sign for $R > 0.1$ and $R < 0.1$ respectively.

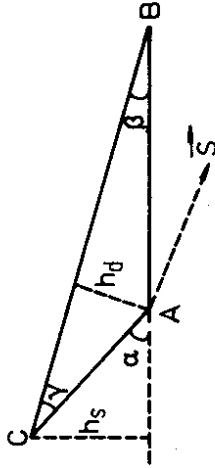


Fig. 2

Fig. 2 The geometric representation of the number of B mesons, N_{Bd} .

The CP nonconservation in the partial decay rates of beauty particles, that does not require mixing at all (e.g., for B^+ or b-baryons) is of interest also. Here the magnitude of the effect cannot be predicted reliably, however, at least for exclusive modes; thus the nontrivial final state interaction (FSI) phases are needed for the asymmetry to appear. For such decays the channels with the interference of double-suppressed $b \rightarrow u\bar{u}s$ transitions with the allowed $b \rightarrow c\bar{c}s$ processes with charmless final states seem to be of particular interest. The CP odd phase difference is α for the above decays.

Now let us briefly discuss the dependence of the number of $B\bar{B}$ events on the unknown parameter R. The required number of $B\bar{B}$ pairs, $N_{B\bar{B}}$, is proportional to

$$N_{B\bar{B}} \sim A^{-2} \cdot B_2 (\beta \rightarrow f)^{-1} \quad (18)$$

The probabilities of the KM-allowed decays $b \rightarrow c$ do not depend on R while the rate of KM-suppressed decays is proportional to

$|V_{ub}|^2 \sim R$. Therefore, the required number of $B\bar{B}$ for both types of the B_d^0 decays is proportional to $|h_d|^{-2}$,

$$(19)$$

$$N_{B_d \bar{B}_d} \sim |h_d|^{-2},$$

where h_d is the perpendicular of the UT, as shown in Fig. 2. Hence $N_{B\bar{B}}$ actually does not depend on R owing to stability of the angle β . For the allowed decays like $B_d(\bar{B}_d) \rightarrow J/\psi K_S$ both the asymmetry and the branching ratios do not change while for KM-suppressed channels like $B_d(\bar{B}_d) \rightarrow \pi^+ \pi^-$ the increase of the asymmetry completely compensates the decrease of probability for smaller values of R.

As for the KM-suppressed decays of B_s , say, $B_s(\bar{B}_s) \rightarrow J/\psi K_S$, the $N_{B\bar{B}_s}$ is connected with the perpendicular h_s of the UT (see Fig. 2). Here the asymmetry increases somewhat slower than the branching ratios decrease, and the required statistics slightly increases when R goes down:

$$N_{B_s \bar{B}_s} \sim \left(\frac{1}{1 + \sqrt{10R}} \right)^2. \quad (20)$$

In fact $N_{B_s \bar{B}_s}$ increases only by a factor of 3 for R varying from 0.2 to 0.015. Thus, the discussed CP odd effects appear to be quite stable with respect to changing the unknown parameter R from the purely statistical point of view [7,15].

Let us recall that all the angles in the UT are roughly proportional to X_{Bd}^{-1} . Therefore the required number of $B\bar{B}$ pairs is rather stable also under the reasonable variation of the experimental value of the $B_d - \bar{B}_d$ mixing parameter X_{Bd} . For B_s mesons, where the rough estimate of the mixing parameter X_{B_s} is $X_{B_d} \geq 7$, only the studies of the time-dependent effects seem to be the effective way to find the CP nonconservation [19,20]. Consequently only the variation of the CP odd phase ϕ affects the required statistics,

$$(21)$$

$$N_{B_s \bar{B}_s} \sim X_{B_d}^2.$$

The dependence of $N_{B\bar{B}}$ on the hadronic matrix elements and $|V_{bc}|^2$ for any B particle follows directly from Eq(11)

$$N_{B\bar{B}} \sim B_K^2 f_B^{-4} \tau_B^{-4}. \quad (22)$$

The discussion above enables one to trace easily the variation of the statistics needed for the detection of the CP nonconservation in the SM under the change of different parameters. Let us turn to the numerical estimates.

Up to now the general strategy for searching the CP violating signals in B decays looks more or less definite only for $e^+ e^-$ collisions. Here there are two distinct alternatives, threshold machines and Z^0 factories.

It appears that to see an effect one needs about $10^8 B\bar{B}$, so the machines with high luminosity $L \sim 10^{33} - 10^{34} \text{ cm}^{-2} \text{ sec}^{-1}$ are required.

In the SM the largest effects are expected for $B_s^0 \bar{B}_s^0$ mesons in the processes with the interference of the decay channels $B^0 \rightarrow f$ and $B^0 \rightarrow \bar{B}^0 \rightarrow f$. However here one should tag the flavour of the initial B-particle, since the final state by itself cannot distinguish between B^0 and \bar{B}^0 . The convenient way is to study the correlation with, say, the sign of the lepton from the semileptonic decay of one of the B's (see e.g. review [16]).

$$\alpha = \frac{N(e^- X + f) - N(e^+ X + \bar{f})}{N(e^- X + f) + N(e^+ X + \bar{f})}. \quad (23)$$

As well known, this asymmetry vanishes when initial $B^0 \bar{B}^0$ pair is born in the C-odd state, for example, in $\Upsilon(4S)$ resonance. The reason is that for the antisymmetric C-odd initial $B^0 \bar{B}^0$ state the interference term \mathcal{M}_{int} is antisymmetric with respect to time intervals t_e and t_f between the birth of $B^0 \bar{B}^0$ pair and the decays to a lepton and to the f state:

$$\mathcal{M}_{\text{int}} \sim e^{-i(t_e + t_f)} \cdot \sin \Delta m (t_f - t_e). \quad (24)$$

Therefore any symmetric treatment of both decay vertices, in particular, the integral measurements, would cancel the effect for the C-odd $B^0 \bar{B}^0$ pairs. To observe the CP odd effects for $B^0 \bar{B}^0$ one needs either the C-even initial states, or to measure the

time-dependence of B decays. For the latter case it is sufficient to determine only the sign of the proper time difference $\Delta t = t_f - t_l$: if one accounts the asymmetry $N(l^+x+f) - N(l^+x+\bar{f})$ with the opposite signs for the decays with $t_f > t_l$ and $t_f < t_l$. Then the cancellation disappears, and the corresponding asymmetry \tilde{a} would coincide with that for the tagged B particles [15].

For f, being a CP eigenstate \tilde{a} is given by

$$\tilde{a} = -2f \sin 2\Phi \frac{x}{1+x^2} \quad (25)$$

(compare with Eq(16)). For realization of such measurements the asymmetric e^+e^- beams (E_e, E_{e^-}) are especially convenient, since there the moving $\Upsilon(4S)$ is produced. Indeed, in the rest frame of $\Upsilon(4S)$ the velocity β_B is only about $\beta_B \approx 0.06$, i.e. for $C\tau_B \approx 300 \mu\text{m}$ the mean decay length $l_B = \beta_B C\tau_B$ is small, $l_B \sim 20 \mu\text{m}$.

On the other hand even for the moderate e^+e^- energy asymmetry, $6.5 \times 4.3 \text{ GeV}$, the velocity β_B is 3.5 times larger. The corresponding distances seem to be sufficient for the vertex detectors with the resolution about $30 \mu\text{m}$ which are now discussed. Moreover, for such collisions the difference of longitudinal coordinates of both vertices is, roughly, proportional to the time difference $\Delta t = t_f - t_l$.

In the Table (see Ref[15] for details) we presented the estimates of branching ratios, the reconstruction efficiencies, the numbers of reconstructed events, N_{reconst} , and the expected statistical significance for some decay channels of B_d mesons. The integral luminosity was assumed to be $L_{\text{int}} = 10^8 \text{ nb}^{-1}$. This corresponds to $N_{B\bar{B}} \approx 10^8$ at $\Upsilon(4S)$. The tagging efficiency ξ_{tag} was taken to be $\xi_{\text{tag}} \approx 0.2$. One can see that the KM suppressed $b \rightarrow u\bar{d}$ transition looks more optimistic as compared with the allowed decays. The reason is the worse reconstruction efficiency for charmed final states and the colour suppression for the particular $b \rightarrow c\bar{c}s$ decays, which takes place in the calculation of relevant branching ratios.

channel	Br [21]	$\mathcal{E} \cdot \text{Br}$	N_{reconst}	A^{CP}/σ_A
$J/\psi K_S$	$3 \cdot 10^{-4}$	$2 \cdot 10^{-5}$	400	1.2
$J/\psi K_S^0 \rightarrow K_S \pi^0$	$4 \cdot 10^{-4}$	$2 \cdot 10^{-5}$	400	$\leq (1.2, 6)$
$J/\psi \pi^+ \pi^-$	10^{-4}	$7 \cdot 10^{-5}$	1400	2.5
$J/\psi \rho^+ \rho^-$	$3 \cdot 10^{-4}$	$15 \cdot 10^{-4}$	3000	3.1
$J/\psi A_1^+ A_1^-$	$3 \cdot 10^{-4}$	10^{-4}	2000	2.6
$\rho^+ \rho^-$	$2 \cdot 10^{-4}$	$8 \cdot 10^{-5}$	1600	$\leq (2.4, 6)$
$\rho^+ A_1^-$	$3 \cdot 10^{-4}$	$8 \cdot 10^{-5}$	1600	$\leq (2.4, 6)$

Let us emphasize that some of the final states above are not CP eigenstates. In this case the effect depends also on the absolute value of the ratio of the $\bar{B}^0 \rightarrow f$ and the $B^0 \rightarrow f$ amplitudes, $|g_f|$, as well as on the CP-even FSI phase difference in both transitions. It is important that the measurement of the time dependences for the decays to f and to \bar{f} enables one to determine $|g_f|$ (and both the CP odd and CP even phases encountered separately; moreover, for B_d mesons, where the mixing is not too large, the extraction of the CP odd phase can be done actually without sizeable loss of statistical significance, at least for the moderate FSI phase difference. The remaining uncertainty in the relative sign of the corresponding amplitudes and connected ambiguity in sign of the CP odd phase can be resolved by measuring the probability of another kind of reactions.

If one of B mesons decays to the CP eigenstate g, then (neglecting the CP nonconservation) the second B meson is really in the state B_1 or B_2 having the CP parity opposite to that of the state g. Indeed, due to Bose-statistics only the $B_1 B_2$ system can be produced in the antisymmetric state with $L = 1$. Then the probability of the decay of the second B meson to the state f is given by $|A(B \rightarrow f) + \eta_g A(\bar{B} \rightarrow f)|^2$, therefore it depends crucially on the relative sign of these two amplitudes.

6) The mixture of different orbital momenta may reduce the asymmetry

Thus with 10^8 $B\bar{B}$ pairs at $\mathcal{Y}(4S)$ the combined effect over the $\pi^+\pi^-\bar{A}_1^+$, $\pi^+\pi^-\rho^+$ channels, could be seen at the level of 4.7 standard deviations. As for the other mechanisms for the partial width difference, which do not imply $B^0-\bar{B}^0$ mixing, say, the interference of the penguin processes with the tree amplitudes, the predictions are far less definite. The crude estimates show that at realistic assumptions about FSI phases such exclusive decays require the order or two higher statistics as compared with the analogous decays with $B^0-\bar{B}^0$ mixing interference. Such effects are also rather complicated for the interpretation in terms of the fundamental parameters of the theory.

The last problem above is not so drastic for the inclusive CP asymmetry in the decays to the strange noncharmed states

$$\delta = \frac{\Gamma(\bar{b} \rightarrow \bar{s}) - \Gamma(b \rightarrow s)}{\Gamma(\bar{b} \rightarrow \bar{s}) + \Gamma(b \rightarrow s)}, \quad (26)$$

arising owing to the interference of penguins with doubly-suppressed $b \rightarrow u\bar{u}d$ transition. Here at $R \approx 0.1$ the asymmetry δ is expected to be about $-5 \cdot 10^{-3}$, with the inclusive rate being near $4 \cdot 10^{-3}$ at $R \approx 0.1$. Thus, purely statistically one needs here about 10^8 $B\bar{B}$ also; however there is a serious background from the allowed processes $b \rightarrow c\bar{u}d$ with two orders of magnitude larger rate.

It is interesting to note that, formally speaking, the required number of B mesons for such inclusive measurements even without substantial suppression of the background decays is just $N_{B\bar{B}} \sim 10^{10}$, that is needed for the inclusive channels like $B^+ \rightarrow K^+\rho^0$.

The real difficulty in searching the CP violating signals in B decays is the necessity to look for the rare few-particle exclusive modes. It is easy to see, for example, that for the decays of neutral B-mesons to the CP eigenstates the sign of the interference is correlated with the final state CP parity.

For this reason the sign of the CP asymmetry is different in the decays proceeding through the same quark transition but leading

to final states with the opposite CP parity. Therefore one may expect the significant decrease of the CP asymmetry when one simply sums up the large number of different events corresponding to the particular quark transition. It is easy to estimate the resulting suppression factor $\langle \eta \rangle$ for the case of "semi-inclusive" measurements when all final states for the given quark decay are included, if one relies on the simple quark diagram approach. It is clear enough that for such decays of B_d mesons as $b \rightarrow u\bar{u}d$ or $b \rightarrow c\bar{c}d$, and for $b \rightarrow c\bar{c}s$ or $b \rightarrow u\bar{u}s$ transitions in B_s -mesons, leading to genuine flavourless states, the factor $\langle \eta \rangle$ is just the average value of the CP parity. Our estimates are as follows:

decays to	B_d	B_s
charmless states	0.03	0.05
one c or \bar{c} quark	0.05	0.09
$c\bar{c}$	0.11	0.19

Note that $\langle \eta \rangle$ appears to be proportional to f_B^2 . We have used quite moderate values $f_{B_d} = 100$ MeV and $f_{B_s} = 130$ MeV as well as not too small phase space suppression factor 0.2 for $b \rightarrow c\bar{c}d(s)$ decays. Therefore even under these conservative assumptions the semi-inclusive measurement in the $b \rightarrow c\bar{c}d$ decays of B_d seems to be comparable with the exclusive decay $B_d^0 \rightarrow \pi^+\pi^-$; the latter exhibits an order of magnitude larger asymmetry with two orders smaller rate.

Note that in the allowed $b \rightarrow c\bar{c}s$ decays of B_d and in the $b \rightarrow u\bar{u}d$ decays of B_s the final states have the opposite strangeness for B and \bar{B} , therefore in the total rates the effect vanishes and $\langle \eta \rangle = 0$. For the decays to the neutral kaons the absence of the asymmetry is due to the cancellation between the states with K_S^0 and K_L^0 . Actually one has not only to separate K_S^0 and K_L^0 to have a sizeable asymmetry here but also to effectively reject the chains like $K^{*+} \rightarrow K_S^0 \pi^+$ which tag strangeness in fact by the sign of the accompanying pion.

Let us return for a moment to the UT. It is known that the large value of m_t leads to the decrease of the theoretical predictions

for the parameter $(\mathcal{E}'/\mathcal{E})_K$ in the SM. This parameter can be estimated as

$$(\mathcal{E}'/\mathcal{E})_K \approx 2.25 \cdot 10^{-2} \cdot \frac{\sqrt{10R}}{1+R} \cdot \sin \Delta, \quad (28)$$

where the numerical factor is proportional to $|V_{cb}|^2$ and to the value of the certain hadronic matrix element, describing the contribution of penguin graphs. One sees that the $(\mathcal{E}'/\mathcal{E})_K$ ratio is actually proportional to the length of the perpendicular h_s' , shown in Fig. 2. Therefore it decreases very slowly with the decreasing of R . The mean value for $R \approx 0.1$, $\mathcal{E}'/\mathcal{E} \approx 3 \cdot 10^{-3}$, appears to be in a good agreement with the recent experimental results. Note that we use the value $H = 0.7$ for the penguin-induced matrix elements, which is needed for the explanation of the $I = 1/2$ rule in $K \rightarrow 2\pi$ decays.

Concluding this section let us emphasize that the e^+e^- threshold machines probably will be the principal source of information on CP violation in B -mesons for the first period of investigations. The long range prospects seem to be connected with the hadron collisions.

In any case the ARGUS results encourage the hopes that in the not so far future CP violation in the beauty hadrons will be observed.

4. B MESON DECAY INTO THE LIGHT HIGGS [22]

As well known, the large value of m_t would enhance significantly a number of the loop-induced weak decay amplitudes of beauty particles. We dwell here on the decay $b \rightarrow s+H$ for sufficiently light Higgs boson H . In the SM the $b \rightarrow s+H$ amplitude is known to be quadratic in m_t leading to an estimate (see e.g. [23])

$$B_2(b \rightarrow s+H) \approx 0.38 \left(\frac{m_t}{80 \text{ GeV}} \right)^4 \left(1 - \frac{m_H^2}{m_b^2} \right)^2. \quad (29)$$

However the existence of very light Higgs boson in the SM looks very inconceivable. The decent motivation for light scalar Higgs could provide SUSY. In the minimal SUSY SM with $m_H^2, m_{H^\pm}^2, m_A^2$ is neutral pseudoscalar) the lightest Higgs has just the standard

tree couplings with fermions and gauge bosons, the tree-induced decays like $J/\psi \rightarrow H + \gamma$ or $\Upsilon \rightarrow H + \gamma$ are not modified. At the same time the radiatively induced amplitudes like $b \rightarrow s+H$ would drastically change due to presence of new particles in loops. It was shown in Ref[22] that in realistic SUSY-SUGRA models with light Higgs boson the typical $b \rightarrow s+H$ amplitude has the opposite sign as compared with the standard one, A_0 .

Its magnitude, being $\sim 4/3 A_0$ in the case of the weak SUSY breaking, while A_0 for strongly broken SUSY, depends on the masses of H^\pm boson and of sparticles. The typical values appeared to be about $(-0.5 - 0.7) A_0$ for $m_t \sim 80$ GeV, with the contribution of the H^\pm exchange being dominant. For lighter top quark, $m_t \sim 50-60$ GeV the SUSY violation becomes stronger and more drastic cancellations generally occur, leading to the values of the $b \rightarrow s+H$ amplitude typically of the order of $(-0.4 + 0.2) A_0$. Thus we believe that the recent data [24] on a search for the Higgs boson in B meson decays do not exclude a light SUSY Higgs boson.

5. CONCLUSIONS

The new experimental results have demonstrated the overwhelming phenomenological successes of the SM. In particular, they allow one to constrain the parameters of this model even in its "hidden sector". The phenomena discussed are very sensitive to the presence of New Physics. The signals for the latter may come from the discovery e.g. of the relatively light top, observation of the not rapid $B_s - \bar{B}_s$ mixing or of the large \mathcal{C}_B , from violation of the "planar trigonometry", from comparison of the m_t with the precise value of M_Z etc.

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