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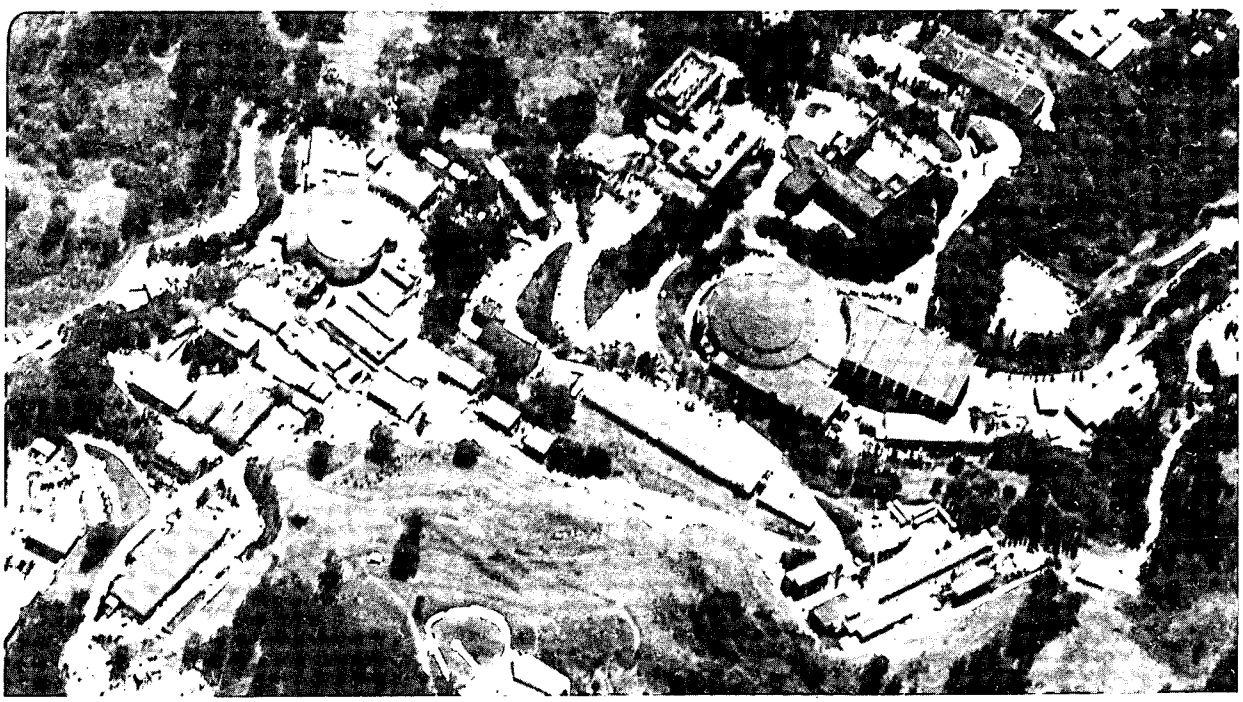

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Aspects of B Physics *

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1. Probing the standard model.

As is well known, the study of B–decays plays an essential role in the determination[1] of the elements of the Cabibbo–Kobayashi–Maskawa (CKM) matrix[2] of charged–current weak couplings, including the CP violating phase δ. In addition it can shed new light on nonleptonic decay dynamics;[3] here the buzz words are factorization, annihilation and penguins. I will comment in more detail below on various aspects of weak decays.

Perturbative QCD can be further probed by, for example, measurements of inclusive $b\bar{b}$ production in hadron colliders. These measurements are of considerable intrinsic interest, particularly the effects of mass dependence on the approach to scaling. They are also important for extrapolating the $b\bar{b}$ content of structure functions to very high energy machines like the SSC or the LHC. They are in any case an essential prerequisite for meaningful decay studies, in particular for CP violating searches, at hadron colliders.

Somewhere on the borderline between perturbative and nonperturbative QCD, there is the possibility for studying the spectroscopy and static properties of a heavy/light bound state system. Interpolation from the presumably relativistic $b\bar{u}, b\bar{d}$ and $b\bar{s}$ systems to the approximately nonrelativistic $b\bar{c}$ system could provide new insight on quark bound states in QCD.

In the realm of nonperturbative QCD, lattice calculations are a growing industry with an ever widening range of applications. Measurements of f_B, the leptonic B decay constant, and the infamous "bag factor" B_B, to be defined below, can be confronted with the results of lattice calculations,[4] as well as of other calculational techniques such as QCD sum rules[5] and the $1/N$ expansion.[6] These parameters play an important role in the analysis of CP violation and the CKM matrix, so reliable information is needed.

Lattice calculations are now also being applied to the determination of structure functions. One could imagine that experimental measurements of the heavy quark content of structure functions or heavy quark fragmentation functions would eventually be able to further test these calculational techniques.

2. Probing beyond the standard model.

One tool for probing new physics is the CKM analysis itself. An inconsistency among data could be interpreted as a signal for a nonstandard effect. An obvious possibility is the existence of one or more additional generations of quarks and leptons. Couplings of known quarks to those of heavier generations destroy the unitarity of the 3×3 CKM submatrix for couplings of the first three generations. Thus at some level a discrepancy should show up.

In addition new particles could mediate additional contributions to decay matrix elements, either at tree level or at the one loop level that determines $B-\bar{B}$ mixing. Candidate particles include additional Higgs bosons, which are expected, in particular, in supersymmetric extensions of the standard models that also predict superpartners for all known particles. Superstring inspired models suggest even more exotic new scalars as well as fermions. New gauge bosons can be present in left–right symmetric extensions of the standard electroweak gauge theory and in some superstring inspired versions.

Another probe of new physics is the study of rare decays.[7] These might involve emission of a new pseudoscalar particle

$$B \to a + X \text{ or } f + X \tag{1}$$

where a is an axion first suggested[8] in the context of a Peccei–Quinn $U(1)$ symmetry invoked[9] to suppress strong CP violation. Axions tend to turn up naturally in supersymmetric models, especially superstring–inspired ones.

In a different vein, some theorists attempt to understand the observed patterns of fermion masses and mixing in terms of a "horizontal symmetry", i.e. a symmetry that interchanges particles of different generations with the same $SU(3)_c \times SU(2)_L \times U(1)$ quantum numbers.[10] This symmetry is obviously broken since these states are not degenerate. If it is a spontaneously broken global symmetry there are necessarily associated goldstone bosons that can be emitted[11] in flavor changing neutral transitions; these are the "familons" f of Eq. (1).

If the horizontal symmetry is gauged, there are neutral gauge bosons that directly mediate lepton and quark flavor–changing interactions, so one expects[10] decays like

$$B \to \left\{ \begin{matrix} \tau\mu \\ \tau e \\ \mu e \end{matrix} \right\} + X. \tag{2}$$

Such decays are also predicted in extended technicolor models[12] that have been constructed in attempts to solve the gauge hierarchy problem. They can also be induced by some of the exotic particles of superstring inspired models.

3. Why B's? (Theory)

Theorists view B–mesons as heavier replicas of K–mesons. The point is that D–decay, as will be T–decay, is dominated by fast CKM allowed transitions: $c \to s$ and $t \to b$. On the other hand K and B decays can proceed only through first forbidden transitions: $s \to u$ and $b \to c$. This means that rare processes have enhanced branching ratios. B decay of course provides an additional probe of the CKM matrix through its second forbidden $b \to u$ transition.

CKM suppression of decay rates also enhances flavor changing $|\Delta F| = 2$ transitions that induce meson–anti–meson mixing and superweak CP violation, since these necessarily entail at least first CKM forbidden couplings. The loop diagrams of Fig. 1 induce mass mixing via a mass difference Δm between eigenstates.[13] The GIM mechanism,[14] i.e. unitarity of the CKM matrix, assures that these diagrams cancel exactly in the limit that the internal quark masses are degenerate. The relatively long B lifetime tells us that the first two generations couple rather weakly to the third. This means that Δm_K nearly vanishes for $m_u = m_c$, obviously a bad approximation relative to the scale m_K, but $\Delta m_D \simeq 0$ for $m_s = m_d$, which is an excellent approximation relative to the scale m_D. For Δm_B, t–exchange is comparable to c, u exchange so the GIM cancellation is again badly broken, but it will be quasi–exact for Δm_T.

The width difference $\Delta\Gamma$ can be schematically represented[15] by the cut diagrams of Fig. 2, i.e. the absorptive parts of Fig. 1. For K and B the GIM cancellation is *a fortiori* broken by the fact that the decay energy is below the threshold for charm and top emission, respectively. However, there is here an essential difference between neutral kaons and B–mesons. K–decay has very limited phase space. As a consequence, approximate CP invariance implies that K_S decays almost exclusively into two pions while K_L has only phase space suppressed 3–body final states: $\Gamma_L << \Gamma_S$. In this case "width mixing" is maximal: $\Delta\Gamma_K/\Gamma_K \sim 1$. In contrast, B decays have a large energy release, so that many channels are open for both CP modes and one expects $\Gamma_1 \simeq \Gamma_2$ for the decay eigenstates. Another difference is that nonleptonic decays of K^0 and $\bar{K}^0$ are into the same (first forbidden) $(n\pi)^0$ channels. Decay channels can be common to B^0 and $\bar{B}^0$ only through an additional Cabibbo suppression factor except for $B_s^0 \leftrightarrow c\bar{c} + X \leftrightarrow \bar{B}_s^0$ which is phase–space inhibited. The net result of these effects is that one expects $\Delta\Gamma_B/\Gamma_B << 1$.

The situation is somewhat different when one considers mass mixing and superweak CP violation. In the standard model, observable CP violation can occur only to the extent that a process probes the existence of all three quark generations. To the extent that the third generation decouples from the first two, the loop diagrams, Fig. 1, that determine $K^0 - \bar{K}^0$ mass mixing are dominated by c and u exchange, so one gets[13]

$$\Delta m_K^2 \propto f(m_c^2) G_F^2 \theta_c^2 \tag{3}$$

where the function[16] $f(m_c^2) \simeq m_c^2$ for $m_c^2 << m_W^2$, so

$$\Delta m_K/\Gamma_K \propto m_c^2/m_K^2 \sim 1 \tag{4}$$

On the other hand CP violation, which determines[17] the imaginary part of Δm_K, requires t–exchange. The smallness of the observed CP-violation in the kaon system can be understood

in terms of the small $s-t$, and very small $d-t$, couplings.

For $B_d - \bar{B}_d$ mixing the presence of the first and third generations in the $(b\bar{d})$ and $(\bar{b}d)$ external states means that the existence of three generations is implicit in the process. More precisely if I denote by $\theta << 1$ the degree of forbiddeness of a transition, and by θ_{ij} the CKM matrix element for $i \leftrightarrow j$, the observed CKM pattern

$$\left.\begin{array}{llllll} \theta_{us} & \equiv \theta_c & \sim \theta_{cd} & \sim \theta_{cb} & \sim \theta_{ts} & \sim \theta \\ \theta_{ub} & \sim \theta_{td} & \sim \theta^2 & & & \end{array}\right\} \tag{5}$$

implies that u, c, and t exchange are of roughly comparable importance for $B_d \leftrightarrow \bar{B}_d$. One finds in fact that the imaginary part of Δm_{B_d} is simply determined by the CP violating phase parameter in the CKM matrix:[18]

$$\arg(\Delta m_{B_d}) \simeq \tan\delta. \tag{6}$$

On the other hand $B_d - \bar{B}_d$ mixing itself is doubly CKM forbidden

$$\Delta m_{B_d}/\Gamma_{B_d} \propto \theta^2 f(m_t^2) \tag{7}$$

and therefore small[19] unless the top quark mass m_t is large.

In contrast Δm_{B_s} is dominated by c and t exchange and suffers no CKM suppression relative to the decay width

$$\Delta m_{B_s}/\Gamma_{B_s} \propto f(m_t^2), \tag{8}$$

but since the first generation is now relatively unimportant, CP violation is expected to be small.

Note that for large m_t, $B - \bar{B}$ mass mixing grows roughly linearly with m_t^2, whereas $\Delta\Gamma_B$ remains fixed (except for the mild, logarithmic dependence of Penguin diagrams – see below – on m_t). Observable "superweak" CP violation arises through a clash of CP violating phases[20] of the contributions of Figs. 1 and 2 and is measurable only if they give comparable contributions to meson–anti–meson mixing. For example in the standard model the charge asymmetry in same sign dilepton events, $B \to \ell^\pm\ell^\pm + X$:

$$A \equiv \frac{\ell^+\ell^+ - \ell^-\ell^-}{\ell^+\ell^+ + \ell^-\ell^-} \tag{9}$$

is, for small $|\Delta\Gamma/\Delta m|$ given by[20]

$$A \simeq \sin\delta\,|\Delta\Gamma/\Delta m|\,. \tag{10}$$

This is expected to be negligibly small for values of m_t as large as those suggested by the data to be analyzed in the next section. This means that one will have to look for CP violating

signals by studying particular decay channels, each of which will have a small branching ratio.

B–decay should also provide an interesting new probe of weak decay dynamics.[3] The collective wisdom is that since B–mesons involve a heavier quark, with a higher energy release in their decays, QCD corrections to weak amplitudes should be smaller and better understood than for the lighter K– and D–mesons.

This wisdom is almost certainly well founded for some applications, notably for penguin diagrams,[21] depicted in Fig. 3. For B–decay, the loop momentum is effectively cut off at the scale $\mu \simeq m_t$. At this scale the effective QCD fine structive constant α_S is small, so the penguin contribution should be well approximated by the leading single gluon exchange diagram which gives an amplitude

$$A_{Penguin} \propto \theta \frac{\alpha_S(m_t^2)}{4\pi} \ln(m_t^2/m_c^2). \tag{11}$$

In contrast, for K–decay the loop momentum is cut off at the much lower scale $\mu \simeq m_c$.

Penguin diagrams in B–decay are particularly interesting because they yield a distinctive final state: $B \to K + \cdots$ at the first CKM forbidden level. Therefore, selecting final states with strange particles that are not charm decay products should enhance the penguin contribution. The competing decay mechanism is a third CKM forbidden transition:

$$\begin{array}{l} b \to u + W^- \\ \qquad\quad \hookrightarrow s\bar{u}. \end{array} \tag{12}$$

Since the matrix elements for (11) and (12) have different phases in the standard model,[2] these final states may also be a good laboratory for studying CP violation. Bjorken[22] has estimated that the two contributions should be of comparable magnitude with the $\alpha_S(m_t^2)$ suppression of the penguin diagrams approximately compensating the extra θ^2 suppression of the decay (12).

QCD corrections also govern the value of the "bag parameter", so named because it was first estimated[23] in the context of the MIT bag model. Specifically, for a neutral pseudoscalar P, $P - \bar{P}$ mass mixing is determined at the quark level by the diagrams of Fig. 1. After loop integration the resulting effective quark–field operator is a $V - A$ current–current operator, shown schematically in Fig. 4a. The matrix element of this effective operator between P and $\bar{P}$ states, which determines Δm_P, is parameterized as

$$\Delta m_P \propto m_P^2 f_P^2 B_P. \tag{13}$$

For $B_P = 1$ this is just determined[13] by the squared P-to-vacuum matrix element

$$\langle P|J_\mu|0\rangle = p_\mu f_P \tag{14}$$

evaluated on the P mass shell: $p^2 = m_P^2$. In the context of QCD, after corrections for hard gluon exchange,[24,25] the parameter $B_P \neq 1$ takes into account soft gluon exchange between the two $V-A$ quark currents, Fig. 4b. Conventional wisdom (which seems to be supported by lattice gauge calculations[4,7]) holds that QCD corrections should be small, so that $B_P \simeq 1$, for heavy quark bound states. The physical grounds for this assumption are questionable because a heavy/light bound state is not really a short-distance system. To tighten the analysis of mass mixing and CP violation, as well as to measure (via Δm) the bag parameter B_B, it is important to have independent measurements of f_P for each pseudoscalar. These measurements also provide tests of nonperturbative QCD calculational techniques.[4-6] Thus one would like to know the partial widths for

$$\left.\begin{matrix} D \\ F \\ B \end{matrix}\right\} \to \tau\nu_\tau. \tag{15}$$

In the important B-decay case, this means measuring very small branching ratios. Bjorken has estimated[26] the branching ratio for $B_c(b\bar{c}) \to \tau\nu_\tau$ at about 1.5%; the decay $B_u(b\bar{u})$ is further suppressed by a double CKM factor θ^2.

Another issue in nonleptonic decay dynamics is the importance of "annihililation" diagrams,[18] Fig. 5, relative to the presumably dominant "spectator" diagram,[27-29] Fig. 6.

For free quarks the annihilation processes of Fig. 5 are helicity suppressed for a $J = 0$ final state with quasi-massless quarks. The argument is identical to that which explains the suppression of $K \to e\nu_e$ relative to $K \to \mu\nu_\mu$. Specifically, the B-decay amplitudes arising from the diagrams of Fig. 5 are determined[18] as

$$A_S \propto G_F f_B m_q \tag{16}$$

where in Eq. (16) m_q is the mass of the heaviest final state quark. In QCD, gluon emission, Fig. 7, can modify this result, since the final state $q\bar{q}$ pair no longer has to be in a $J = 0$ configuration.[30] The fact that the $D^+(c\bar{d})$ decays more slowly than the $D^0(c\bar{u})$ or the $F^+(c\bar{s})$ is generally attributed to the presence of the annihilation mechanism. The diagrams of Fig. 7b and Fig. 7a contribute respectively to CKM allowed transitions for D^0 and $F^\pm$. No CKM allowed annihilation process can contribute to $D^\pm$ decay. In contrast, the spectator decay of Fig. 6 is independent of the flavor of the quark bound to the decaying charmed quark. If only this process is important one expects[27-29] equal life-times for D^0, D^+ and F^+, up

to interference effects[31] that may occur in $D^\pm$ decays, due to the fact that the spectator quark is the same as one of the decay products in the CKM allowed transition. In fact, the difference in lifetimes is much less dramatic than indicated by early experiments and could possibly be attributed[32] solely to destructive interference of spectator diagrams in $D^\pm$ decay.

In any case, the common wisdom[33] is that the annihilation process of Fig. 7 should be less important in B-decay because of the higher mass scalar and thus the smaller effective QCD coupling constant. However, as I noted above, a heavy (Q)-light (q) bound quark system is not really a short-distance system. Its inverse size is determined by the reduced mass, which is simply m_q if $m_q << m_Q$. Put another way, the distance over which the annihilation processes of Fig. 7 take place is the requisite "off-mass-shellness" of the virtual quark which, for emission of a massless gluon, is of the order of the mass of the quark from which the gluon is emitted. This suggests that the relevant QCD fine structure constant is $\alpha_S(m_q)$ rather than (the much smaller) $\alpha_S(m_Q)$. Within this perspective the $B_c(b\bar{c})$ system is really "small", in that the relevant distance scale is m_c^{-1}. It would therefore be extremely interesting to study the annihilation process in B-decay as a function of the mass of the lighter bound quark. For the $B_s(b\bar{s})$ system CKM allowed annihilation (Fig. 6a or 7b) is signed by a $c\bar{c}$ final state, e.g.:

$$B_s \to \left\{\begin{matrix} D\bar{D} & + & \pi's \\ \psi & + & \pi's. \\ \eta_c & + & \pi's \end{matrix}\right. \tag{17}$$

Since these final states are also CP eigenstates they may prove useful in the search for CP violating signals.[34] However, each exclusive decay mode of this type is expected [22,26] to have a branching ratio of less than a per cent.

To test the dependence of annihilation diagrams on the mass of the lighter bound quark, the importance of the final states (17) in B_s decay should be compared with the partial life-times for B_c decays to final states accessible via CKM allowed annihilation, namely (Fig. 6b, 7a):

$$B_c \to \left\{\begin{matrix} pions \\ DK + X \end{matrix}\right. . \tag{18}$$

4. Why B's? (Experiment)

The most recent datum supporting the assertion that B-decays are important is that a single experimental measurement, namely of $B_d - \bar{B}_d$ mixing, instantaneously generated a large number of (for the most part good) theoretical papers.

In fact there have been three recent experimental measurements of prime importance for probing the standard model, namely:

1. The observation[35] of B–decay into a noncharmed final state ($p\bar{p} + X$). This (almost certainly) demonstrates the existence of a direct $b\bar{u}W^+$ coupling, without which CP violation would be inexplicable in the framework of the standard model.

2. The observation[36] that $\epsilon'/\epsilon \neq 0$. This is the first positive indication for CP violation other than in $K^0 - \bar{K}^0$ (superweak) mass mixing, and is an equally important result for substantiating the standard model.

3. A substantial $B_d - \bar{B}_d$ mixing.[37] This result was at first sight surprising because $B_d - \bar{B}_d$ mixing was predicted[19] to be rather smaller than observed, under the assumption of a relatively light ($m_t \lesssim 40$ GeV) top quark. However, as I shall outline below, the observed values of ϵ'/ϵ and $B_d - B_{\bar{d}}$ mixing are quite consistent with the standard model, provided that the top quark mass is rather large.[38–40]

The logic of the analysis of these recent results is as follows. The relative yield of same–sign dileptons in $B_d - \bar{B}_d$ events:

$$r_d = \frac{\ell^+\ell^+ + \ell^-\ell^-}{\ell^+\ell^-} \tag{19}$$

determines (neglecting, as argued above, "width mixing", $\Delta\Gamma \simeq 0$) the B mass–mixing parameter which is governed by the diagrams of Fig. 1. For a large top quark mass the dominant contribution is double t–exchange, giving a contribution:

$$\Delta m_{B_d} \propto \theta_{td}^2 f(m_t^2). \tag{20}$$

Using the standard parametrization[1,2] of the KM matrix in terms of three angles $\theta_1 \simeq \theta_c, \theta_2$ and θ_3 and a phase δ, the $t \leftrightarrow d$ matrix element

$$\theta_{td} \simeq s_c s_2 \tag{21}$$

(here $s_i \equiv \sin\theta_i$) is related by unitarity[2,14] of the CKM matrix to other measured CKM matrix elements. The B-decay life–time, dominated by the $b \to c$ transition, determines[18,41] the element

$$|\theta_{bc}| = \left|s_3 + s_2 e^{i\delta}\right|. \tag{22}$$

The CKM suppressed $b \to u$ transition is experimentally bounded. The experimental limit[42] on the branching ratio

$$R = \frac{b \to u}{b \to c} \tag{23}$$

can be, together with B–lifetime measurements,[43] interpreted as a limit on the $b \to u$ transition matrix element:

$$|\theta_{bu}| = |s_c s_3|. \tag{24}$$

The Cabibbo angle, or s_c, is a well measured quantity. The experimental bounds[42] on the ratio R, Eq. (23), imply a small value for s_3. This in turn, together with a rather long[43] B–lifetime that bounds θ_{bc}, Eq. (22), implies that s_2, and hence θ_{td}, Eq. (21), cannot be very large. As a result, the substantial value observed[37] for $B_d - \bar{B}_b$ mixing implies that the function[16] $f(m_t^2)$, which, for $m_t^2 < m_W^2$, grows with m_t^2, must be large. Numerical analyses have been performed by several groups[38–40,44] who for the most part conclude[38–40] that existing data imply at least $m_t \gtrsim 50$ GeV and, more probably, $m_t \gtrsim 100$ GeV. A dissenting view has been registered by one group[44] that claims that present data allow m_t as low as the roughly 20 GeV limit imposed by the nonobservation of $t\bar{t}$ production at PEP and PETRA. However, the latter authors allow values for the unknown parameters in the analysis, namely f_B, B_B and the ratio R of Eq. (23), that most theorists would probably consider as unreasonable. I emphasize once again the importance of independent measurements of these parameters.

Once the observed $B_d - \bar{B}_d$ mixing has been assimilated within the standard model, the resulting restrictions on allowed values for the parameters of this model have implications for other measurable quantities. Consider first CP violation in the $K^0 - \bar{K}^0$ system. Superweak CP violation (i.e. CP violation in mass mixing) is determined[17] by the imaginary parts of the diagrams of Fig. 1. The CP violating part of these diagrams involves t–quark exchange; this contribution grows as m_t^2 for $m_t < m_W$. Thus the parameter ϵ, which measures the CP violating component of Δm_K, grows roughly as m_t^2. On the other hand, direct CP violation in decays other than the 2π isospin zero mode used to define the superweak CP violating phase, and in particular the parameter ϵ' which measures the CP violating phase in the 2π $I = 2$ mode relative to the $1 = 0$ phase, is governed[45] by Penguin diagrams, as in Fig. 3, that grow only logarithmically with m_t. Thus the ratio ϵ'/ϵ decreases[46] with increasing m_t for $m_t < m_W$. As a consequence, the large value for m_t inferred from the $B_d - \bar{B}_d$ mixing measurement[37] is consistent[38] with the small observed[36] value for ϵ'/ϵ in the context of the standard model.

This picture implies predictions for as yet unmeasured quantities. For example it is inferred[38,39] that $B_s - \bar{B}_s$ mixing should be nearly maximal. In addition the $K^+ \to \pi + \bar{\nu}\nu$ branching ratio prediction is sharpened. This occurs[13,47] in the standard model through the loop diagram of Fig. 8, and, since it is GIM suppressed, grows in importance for large m_t. One finds[38]

$$B(K^+ \to \pi^+\nu\bar{\nu}) \simeq (1-8)10^{-10} \tag{25}$$

for $m_t = (50 - 200)$ GeV. (The experimental bound on the parameter $\rho - 1$ where $\rho = m_W/m_Z \cos\theta_w$ implies[48] an upper limit of about 200 GeV on m_t.) It has also been pointed

out[49] that a top quark mass as large as 200 GeV could give a possibly observable branching ratio for the rare neutral current flavor changing B-decay:

$$B(B^+ \to K^+ + \ell^+\ell^-) \sim 10^{-5}. \tag{26}$$

More generally, the measured $B_d - \bar{B}_d$ mixing[37] tightens the values of the CKM matrix elements and hence the quark mass matrices, which, when diagonized, determine the CKM matrix. A recent analysis[50] has all but ruled out specific conjectures for the form of the quark mass matrices, partially based on GUTs models.

The implications of this measurement[37] for physics beyond the standard model has also been analysed.[39] Contributions from supersymmetric partners of ordinary particles and/or additional Higgs scalars are found to enhance $B - \bar{B}$ mixing. Thus a smaller value of m_t than that inferred in the standard model would be compatible with the data if either of these effects are present.[39] On the other hand $B\bar{B}$ mixing is found to be relatively suppressed in left–right–symmetric extensions of the electroweak gauge theory; an even higher top quark mass limit is inferred in the context of these theories.[39]

5. A superstring–inspired example of exotic physics.

Models inspired by the $E_8 \times E_8$ heterotic[51] superstring[52] theory end up in four dimensions[53] with an (already broken) $E_6 \times E_6$ gauge theory. Here E_6 (or a subgroup thereof) describes a pure supersymmetric Yang Mills theory of a so–called "hidden sector" that interacts only gravitationally with observed matter, and E_6 is the GUT of the observed world. The unbroken subgroup of E_6 at scales just below the compactification scale must contain the observed gauge group $SU(3)_c \times SU(2)_L \times U(1)$. Each matter generation fills a 27–plet of E_6 which decomposes under $SU(5)$ as:

$$27 = (\bar{5} + 10) + (5 + \bar{5}) + 1 + 1. \tag{27}$$

In (27) the $(\bar{5}+10)$ supermultiplets contain quarks (q) and leptons (ℓ) and their superpartners, squarks ($\tilde{q}$) and sleptons ($\tilde{\ell}$). Each $(5+\bar{5})$ supermultiplet contains a Higgs (H) and Higgsino ($\widetilde{H}$) super–multiplet that is a weak isopin doublet, as well as a color triplet supermultiplet $(D, \widetilde{D})$ which has the same flavor quantum numbers under $SU(3)_c \times SU(2)_L \times U(1)$ as the right–handed d–quark. There are as many $(5 + \bar{5})$ supermultiplets as matter generations. This means that there is a large number of physical Higgs particles as well as other exotic states. If there are no discreet symmetries to forbid them, there will be generation mixing couplings among the $(\bar{5} + 10)$ and $(5+\bar{5})$ multiplets, which, if the masses of the later are not very large, will induce[54] effective flavor changing neutral current (FCNC) transitions among light particles, via the diagrams of Figs. 9(a-c). In addition, the possibility of d-D mixing

potentially spoils[54] the GIM mechanism that in the standard model forbids the tree level FCNC process of Fig. 9d, where Z' represents on additional neutral gauge boson that is present if the surviving gauge group in four dimensions is larger than the standard one. If present, all of the processes of Fig. 9 would contribute to Δm_B (and Δm_K), and therefore to the parameter r_d of Eq. (19). Neglecting $\Delta\Gamma_B$, the experimental result[37] implies a bound[54]

$$r_d \simeq \frac{1}{2}\frac{\Delta m_B}{\Gamma_B} \lesssim 0.3 \tag{28}$$

or

$$\Delta m_B \lesssim 4 \times 10^{-13}\text{GeV} \tag{29}$$

which is in the ballpark of

$$\Delta m_K = 3.5 \times 10^{-15} GeV. \tag{30}$$

This means that constraints on new phenomena from Δm_B are comparable to those from Δm_K. For example, assuming $m_{H,\tilde{H}} \simeq 100$ GeV, $m_{D,\tilde{D}} \simeq 300$ GeV $\simeq m_{Z'}$ the bounds[54] on new couplings λ involving external b–quarks in the diagrams of Fig. 9:

$$\lambda_b < 10^{-4} - 0.1 \tag{31}$$

are comparable to these involving external s–quarks:

$$\lambda_s < 10^{-5} - 0.1. \tag{32}$$

Note moreover that (31) and (32) are independent, since couplings involving different matter generations are *a priori* independent.

The new couplings suggested by superstring–inspired theories should also induce FCNC semi–leptonic decays. For example the experimental branching ratio bound

$$B(B \to \ell^+\ell^- + X) < 6 \times 10^{-3} \tag{33}$$

implies[54] the limit

$$\lambda_b \lambda_\ell < (0.06)^2$$

on the couplings of b–quarks and leptons to additional Higgs bosons with $m_H \simeq 100$ GeV, Fig. 10, as suggested by some superstring–inspired models.

6. Conclusions

I hope that I have made it clear that any data on B–decays is at present extremely interesting, in that it provides powerful new constraints in analyses of the standard model and extensions thereof.

Thinking about future detectors and/or facilities for B-meson studies should have as the primary objective the ability to study CP violation. This will be difficult. Bjorken[55] has estimated that at least 3×10^7 $B\bar{B}$ production events are needed for meaningful CP violation studies. This is actually his optimistic estimate, revised upward because of the observed substantial $B_d - \bar{B}_d$ mixing that may facilitate observation of CP violating effects in neutral B-decays. Bjorken's reasoning[55] is as follows:

a) A specific state must be reconstructed. This involves either a CKM-forbidden non-charmed final state, a somewhat phase-space suppressed $c\bar{c}$ final state, or a decay chain $B \to D + f$, $D \to f'$ entailing the product of two small branching ratios for fixed f and f'. Therefore, the overall branching ratio for any given final state will be no larger than $10^{-3} - 10^{-4}$.

b) The associated B or $\bar{B}$ must be flavor-tagged by identifying the charge of a decay lepton and/or the strangeness of the hadronic decay products. This will entail another suppression factor of at least 10^{-1}.

c) Sufficient statistics, at least 10^3 events, must be accumulated for a meaningful search for CP violation in a particular channel.

A necessarily prerequisite for CP violation studies is a good knowledge of production rates and distributions and decay branching ratios. Production and decay branching ratios will provide important data for the standard model, as well as sharpen the choices for the best line of attack on CP violation.

A secondary goal for new facilities or detectors is to push as far as possible limits on rare decays. These can provide powerful constraints on proposed extensions of the standard model – or perhaps one day provide a real signal for new physics.

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Figure Captions

1. One loop quark diagrams that contribute to neutral meson–anti–meson mass mixing.
2. The absorptive parts of the diagrams of Fig. 1 that contribute to width differences in neutral $(q\bar{q}')$ meson systems.
3. Penguin diagrams for nonleptonic B–decays.
4. Schematic representations of a) the factorization approximation to the meson–anti–meson matrix element of the effective $\Delta S = 2$ quark operator generated by the diagrams of Fig. 1, and b) QCD corrections to factorization that generate a "bag factor" $B_P \neq 1$.
5. Annihilation diagrams for nonleptonic B–decays in the free quark approximation.
6. Spectator diagram for nonleptonic B–decays.
7. QCD corrections to the diagrams of Fig. 5.
8. One loop contribution to the $s \to d\nu\bar{\nu}$ transition in the standard model.
9. Diagrams that can generate $B_d^0 - \bar{B}_d^0$ and $K^0 - \bar{K}^0$ mixing in superstring–inspired models.
10. Diagram that can generate $B \to X + \ell^+\ell^-$ or $\mu^\pm e^\mp$ in superstring inspired models.

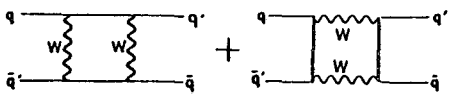

Figure 1

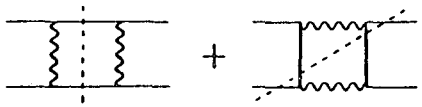

Figure 2

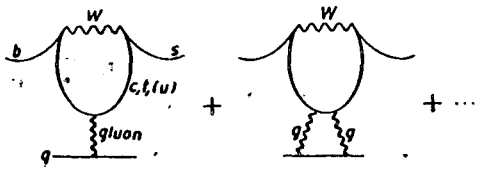

Figure 3

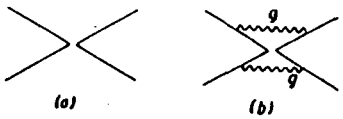

Figure 4

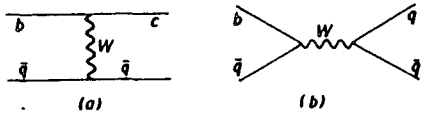

Figure 5

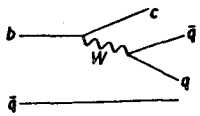

Figure 6

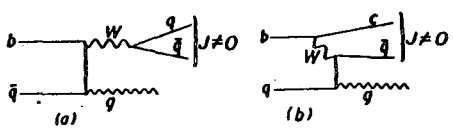

Figure 7

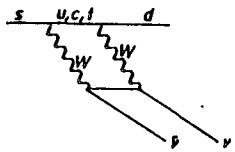

Figure 8

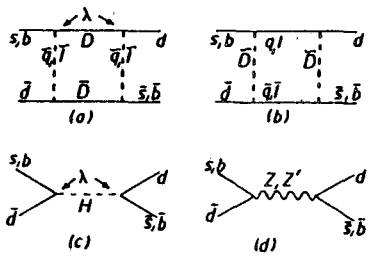

Figure 9

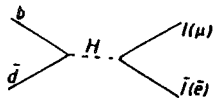

Figure 10

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