PROCEEDINGS

International Europhysics Conference on HEP

New Results on Flavor Physics

Zoltan Ligeti

Ernest Orlando Lawrence Berkeley National Laboratory
University of California, Berkeley, CA 94720
E-mail: zligeti@lbl.gov

Abstract: Recent progress in flavor physics is discussed. In particular, I review theoretical and experimental developments relevant for semileptonic B decays and the determination of $|V_{cb}|$ and $|V_{ub}|$, for exclusive rare decays, for nonleptonic $b \to c$ decays and tests of factorization, and for D meson mixing.

PrHEP hep2001

1. Introduction

The goal of the B physics program is to precisely test the flavor structure of the standard model (SM), that is the Cabibbo-Kobayashi-Maskawa (CKM) description of quark mixing and CP violation. In the last decade the accuracy with which we know that gauge interactions are described by the SM improved by an order of magnitude, and sometimes more. In the coming years tests of the flavor sector of the SM and our ability to probe for flavor physics and CP violation beyond the SM will improve in a similar manner.

However, in contrast to the hierarchy problem of electroweak symmetry breaking, there is no similarly robust argument that new flavor physics must appear near the electroweak scale. Nevertheless, the flavor sector provides severe constraints for model building, and many extensions of the SM do involve new flavor physics which may be observable at the B factories. Flavor physics also played an important role in the development of the SM: (i) the smallness of $K^0 - \overline{K}^0$ mixing led to the GIM mechanism and a calculation of the charm mass before it was discovered; (ii) CP violation led to the proposal that there should be three generations before any third generation fermions were discovered; and (iii) the large $B^0 - \overline{B}^0$ mixing was the first evidence for a very large top quark mass.

The B meson system has several features which makes it well-suited to study flavor physics and CP violation. Because the top quark in loop diagrams is neither GIM nor CKM suppressed, large CP violating effects are possible, some of which have clean interpretations. For the same reason, a variety of rare decays are expected to have large enough branching fractions to allow for detailed studies. Finally, some of the hadronic physics can be understood model independently because $m_b \gg \Lambda_{\rm QCD}$.

In the standard model all flavor changing processes are mediated by charged current weak interactions, whose couplings to the six quarks are given by a three-by-three unitary matrix, the Cabibbo-Kobayashi-Maskawa (CKM) matrix. It has a hierarchical structure, which is well exhibited in the Wolfenstein parameterization,

$$V_{\rm CKM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} = \begin{pmatrix} 1-\frac{1}{2}\lambda^2 & \lambda & A\lambda^3(\rho-i\eta) \\ -\lambda & 1-\frac{1}{2}\lambda^2 & A\lambda^2 \\ A\lambda^3(1-\rho-i\eta) & -A\lambda^2 & 1 \end{pmatrix} + \ldots\,. \tag{1.1}$$

This form is valid to order λ^4. The small parameter is chosen as the sine of the Cabibbo angle, $\lambda \simeq 0.22$, while A, ρ, and η are order unity. In the SM, the only source of CP violation in flavor physics is the phase of the CKM matrix, parameterized by η.

The unitarity of $V_{\rm CKM}$ implies that its nine complex elements must satisfy $\sum_k V_{ik}V_{jk}^* = \sum_k V_{ki}V_{kj}^* = \delta_{ij}$. The vanishing of the product of the first and third columns provides a simple and useful way to visualize these constraints,

$$V_{ud}\,V_{ub}^* + V_{cd}\,V_{cb}^* + V_{td}\,V_{tb}^* = 0\,, \tag{1.2}$$

which can be represented as a triangle (see Fig. 1). Making overconstraining measurements of the sides and angles of this unitarity triangle is one of the best ways to look for new physics.

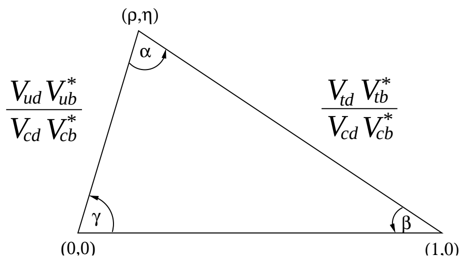

Figure 1: The unitarity triangle

To believe at some point in the future that a discrepancy is a sign of new physics, model independent predictions are essential. Results which depend on modeling nonperturbative strong interaction effects will never disprove the Standard Model. Most model independent predictions are of the form

$$\text{Quantity of interest} = (\text{calculable factor}) \times \left[1 + \sum_k (\text{small parameters})^k\right], \tag{1.3}$$

where the small parameter can be $m_s/\Lambda_{\chi SB}$, $\Lambda_{\rm QCD}/m_b$, $\alpha_s(m_b)$, etc. Still, in most cases, there are theoretical uncertainties suppressed by some (small parameter)^N, which may be hard to estimate model independently. If one's goal is to test the Standard Model, one must assign sizable uncertainties to such "small" corrections not known from first principles.

Over the last decade, most of the theoretical progress in understanding B decays utilized that m_b is much larger than $\Lambda_{\rm QCD}$. However, depending on the process under consideration, the relevant hadronic scale may or may not be much smaller than m_b (and, especially, m_c). For example, f_π, m_ρ, and m_K^2/m_s are all of order $\Lambda_{\rm QCD}$, but their numerical values span more than an order of magnitude. In many cases, as it will become clear below, experimental guidance is needed to decide how well the theory works in different cases.

To overconstrain the unitarity triangle, there are two very important "clean" measurements which will reach precisions at the few, or maybe even one, percent level. One is $\sin 2\beta$ from the CP asymmetry in $B \to J/\psi\, K_S$, which is rapidly becoming the most precisely known ingredient of the unitarity triangle [1]. The other is $|V_{td}/V_{ts}|$ from the ratio of the neutral meson mass differences, $\Delta m_d/\Delta m_s$. The LEP/SLD/CDF combined limit is presently [2]

$$\Delta m_s > 14.6\,{\rm ps} \quad (95\%\ {\rm CL})\,. \tag{1.4}$$

Probably B_s mixing will be discovered at the Tevatron, and soon thereafter the experimental error of Δm_s is expected to be below the 1% level [3]. The uncertainty of $|V_{td}/V_{ts}|$ will then be dominated by the error of $\xi \equiv (f_{B_s}/f_{B_d})\sqrt{B_{B_s}/B_{B_d}}$. For the last few years the lattice QCD averages have been about $\xi = 1.15 \pm 0.06$ [4], surprisingly consistent with the chiral log calculation, $\xi^2 \sim 1.3$ [5]. This year we are learning that an additional error, estimated to be $^{+0.07}_{-0.}$ [4], may have to be added to ξ for now, since in the unquenched calculation chiral logs are important in the chiral extrapolation for f_B, but they do not affect f_{B_s} [6]. It is very important to reduce this uncertainty, and do simulations with three light flavors.

Compared to $\sin 2\beta$ and $|V_{td}/V_{ts}|$, for which both the theory and the experiment are tractable, much harder is the determination of another side or another angle, such as $|V_{ub}|$, or α, or γ ($|V_{cb}|$ is also "easy" by these criteria). However, our ability to test the CKM hypothesis in B decays will depend on a third best measurement besides $\sin 2\beta$ and x_s (and on "null observables"). The accuracy of these measurements will determine the sensitivity to new physics, and the precision with which the SM is tested. It does not matter whether it is a side or an angle. What is important is which measurements can be made that have clean theoretical interpretations for the short distance physics we are after.

Section 2 reviews recent progress for semileptonic decays and the determination of $|V_{cb}|$ and $|V_{ub}|$. Related developments relevant for exclusive rare decays are also discussed. Section 3 deals with nonleptonic decays, such as lifetimes, tests of factorization for exclusive nonleptonic decay, and D^0 mixing. Section 4 contains our conclusions. While this write-up follows closely the slides at the Conference, I attempted to update the experimental results where available. I was asked not to talk about CP violation, which was reviewed in Refs. [1, 7].

2. Semileptonic decays

The determination of $|V_{cb}|$ and $|V_{ub}|$ are very important for testing the CKM hypothesis. The allowed range of $\sin 2\beta$ in the SM depends strongly on the uncertainty of $|V_{ub}|$ (since it determines the side of the unitarity triangle opposite to the angle β), and the constraint from the $K^0 - \overline{K}^0$ mixing parameter ϵ_K is proportional to $|V_{cb}|^4$. This is illustrated in Fig. 2. Moreover, the methods developed to extract $|V_{cb}|$ and $|V_{ub}|$ are also useful for reducing the hadronic uncertainties in rare decays.

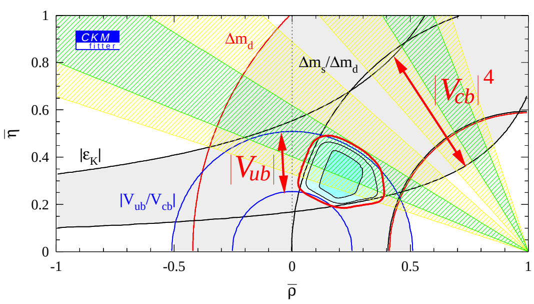

Figure 2: The allowed range of $\bar\rho - \bar\eta$ [8].

2.1 Exclusive $B \to D^{(*)}\ell\bar\nu$ decay and the HQET

In heavy mesons composed of a heavy quark, Q, and a light antiquark, $\bar q$, and gluons and $q\bar q$ pairs, in the $m_Q \to \infty$ limit the heavy quark acts as a static color source with fixed four-velocity v^μ. The wave-function of the light degrees of freedom become insensitive to the spin and mass (flavor) of the heavy quark, resulting in heavy quark spin-flavor symmetries [9].

The determination of $|V_{cb}|$ from exclusive $B \to D^{(*)}\ell\bar\nu$ decays is based on the fact that heavy quark symmetry relates the form factors which occur in these decays to the Isgur-Wise function, whose value is known at zero recoil in the infinite mass limit. The symmetry breaking corrections can be organized in a simultaneous expansion in $\alpha_s(m_Q)$ and $\Lambda_{\rm QCD}/m_Q$ (where $Q = c, b$). The $B \to D^{(*)}\ell\bar\nu$ rates can be schematically written as

$$\frac{{\rm d}\Gamma(B \to D^{(*)}\ell\bar\nu)}{{\rm d}w} = (\text{known factors})\, |V_{cb}|^2 \begin{cases} (w^2-1)^{1/2}\, \mathcal{F}_*^2(w)\,, & \text{for } B \to D^*, \\ (w^2-1)^{3/2}\, \mathcal{F}^2(w)\,, & \text{for } B \to D, \end{cases} \tag{2.1}$$

where $w = (m_B^2 + m_{D^{(*)}}^2 - q^2)/(2m_B m_{D^{(*)}})$. Both $\mathcal{F}(w)$ and $\mathcal{F}_*(w)$ are equal to the Isgur-Wise function in the $m_Q \to \infty$ limit, and in particular $\mathcal{F}(1) = \mathcal{F}_*(1) = 1$, allowing for a model independent determination of $|V_{cb}|$. The zero recoil limits of $\mathcal{F}_{(*)}(w)$ are of the form

$$\mathcal{F}_*(1) = 1 + c_A(\alpha_s) + \frac{0}{m_Q} + \frac{(\ldots)}{m_Q^2} + \ldots\,, \qquad \mathcal{F}(1) = 1 + c_V(\alpha_s) + \frac{(\ldots)}{m_Q} + \frac{(\ldots)}{m_Q^2} + \ldots\,. \tag{2.2}$$

The perturbative corrections, $c_A = -0.04$ and $c_V = 0.02$, have been computed to order α_s^2 [10], and the unknown higher order corrections should be below the 1% level. The order $\Lambda_{\rm QCD}/m_Q$ correction to $\mathcal{F}_*(1)$ vanishes due to Luke's theorem [11]. The terms indicated by $(\ldots)$ in Eqs. (2.2) are only known using phenomenological models or quenched lattice QCD at present. This is why the determination of $|V_{cb}|$ from $B \to D^*\ell\bar\nu$ is theoretically more reliable for now than that from $B \to D\ell\bar\nu$, although both QCD sum rules [12] and quenched lattice QCD [13] suggest that the order $\Lambda_{\rm QCD}/m_Q$ correction to $\mathcal{F}(1)$ is small. Due to the extra $w^2 - 1$ helicity suppression near zero recoil, $B \to D\ell\bar\nu$ is also harder experimentally.

$|V_{cb}|\, \mathcal{F}_*(1)$ is measured from the zero recoil limit of the decay rate, and the results are shown in Table 1. The main theoretical uncertainties in such a determination of $|V_{cb}|$ come from the value of $\mathcal{F}_{(*)}(w)$ at $w = 1$ and from its shape used to fit the data. In my opinion, a reasonable estimate at present is

$$\mathcal{F}_*(1) = 0.91 \pm 0.04\,, \tag{2.3}$$

$\lvert V_{cb}\rvert\, \mathcal{F}_*(1) \times 10^3$	Experiment
35.6 ± 1.7	LEP [14]
42.2 ± 2.2	CLEO [15]
36.2 ± 2.3	BELLE [16]

Table 1: Measurements of $|V_{cb}|\, \mathcal{F}_*(1)$.

where the error can probably only be reduced by unquenched lattice calculations in the future. The quenched result is $\mathcal{F}_*(1) = 0.913^{+0.024+0.017}_{-0.017-0.030}$ [17]. It will be interesting to see the effect of unquenching, and if $|V_{cb}|$ obtained from $B \to D\ell\bar\nu$ using $\mathcal{F}(1)$ from the lattice will agree at the few percent level.

For the shape of $\mathcal{F}_*(w)$, it is customary to expand about zero recoil and write $\mathcal{F}_*(w) = \mathcal{F}_*(1)\, [1 - \rho^2\, (w-1) + c\, (w-1)^2 + \ldots]$. Analyticity imposes stringent constraints between the slope, ρ^2, and curvature, c, at zero recoil [18, 19], which is already used to fit the data and obtain the results in Table 1. Recently there has been renewed effort in constraining the slope parameter ρ^2 using sum rules and data on B decays to excited D states [20, 21]. Decays to orbitally excited D mesons can also be studied in HQET [22, 23], and it seems problematic to accommodate the data which suggests that the rate to the D_1^* and D_0^* states ($s_l^{\pi_l} = \frac{1}{2}^+$) is larger than that to D_1 and D_2^* ($s_l^{\pi_l} = \frac{3}{2}^+$) [20, 21, 23].

Measuring the $B \to D\ell\bar\nu$ rate [24, 25] is also important, since computing $\mathcal{F}(1)$ on the lattice is no harder that $\mathcal{F}_*(1)$, and so it provides an independent determination of $|V_{cb}|$. Comparing the shapes of the $B \to D^*$ and $B \to D$ spectra may also help, since it gives additional constraints on ρ^2, and the correlation between ρ^2 and the extracted value of $|V_{cb}|\,\mathcal{F}_*(1)$ is very large [26].

2.2 Inclusive semileptonic B decay and the OPE

Inclusive B decay rates can be computed model independently in a series in $\Lambda_{\rm QCD}/m_b$ and $\alpha_s(m_b)$, using an operator product expansion (OPE) [27]. The results can be schematically written as

$$\mathrm{d}\Gamma = \begin{pmatrix} b \text{ quark} \\ \text{decay} \end{pmatrix} \times \left\{ 1 + \frac{0}{m_b} + \frac{f(\lambda_1, \lambda_2)}{m_B^2} + \ldots + \alpha_s(\ldots) + \alpha_s^2(\ldots) + \ldots \right\}. \tag{2.4}$$

The $m_b \to \infty$ limit is given by b quark decay, and the leading nonperturbative corrections suppressed by $\Lambda_{\rm QCD}^2/m_b^2$ are parameterized by two hadronic matrix elements, usually denoted by λ_1 and λ_2. The value $\lambda_2 \simeq 0.12\,\mathrm{GeV}^2$ is known from the $B^* - B$ mass splitting. At order $\Lambda_{\rm QCD}^3/m_b^3$ seven new and unknown hadronic matrix elements enter, and usually naive dimensional analysis is used to estimate their size and the related uncertainty. For most quantities of interest, the perturbation series are known including the α_s and $\alpha_s^2\beta_0$ terms, where $\beta_0 = 11 - 2n_f/3$ is the first coefficient of the β-function (which is large, so in many cases this term is expected to dominate the α_s^2 corrections).

The good news from the above is that "sufficiently inclusive" quantities, such as the total semileptonic width relevant for the determination of $|V_{cb}|$, can be computed with errors at the $\lesssim 5\%$ level. In such cases the theoretical uncertainty is controlled dominantly by the error of a short distance b quark mass (whatever way it is defined). Using the "upsilon expansion" [28] the relation between the inclusive semileptonic rate and $|V_{cb}|$ is

$$|V_{cb}| = (41.9 \pm 0.8_{(\rm pert)} \pm 0.5_{(m_b)} \pm 0.7_{(\lambda_1)}) \times 10^{-3} \left(\frac{\mathcal{B}(\bar B \to X_c\ell\bar\nu)}{0.105} \, \frac{1.6\,\mathrm{ps}}{\tau_B} \right)^{1/2}. \tag{2.5}$$

The first error is from the uncertainty in the perturbation series, the second one from the b quark mass, $m_b^{1S} = 4.73 \pm 0.05\,\mathrm{GeV}$ (a conservative range of m_b may be larger [29]), and the third one from $\lambda_1 = -0.25 \pm 0.25\,\mathrm{GeV}^2$. This result is in agreement with Ref. [30], where the central value is 40.8×10^{-3} (including the 1.007 electromagnetic radiative correction).

LEP and BELLE reported new results for the semileptonic branching ratio, which yield a determination of $|V_{cb}|$ which is dominated by theoretical errors,

$$\mathcal{B}(B \to X\ell\bar\nu) = \begin{cases} 10.65 \pm 0.23\% & \text{(LEP [14])} \\ 10.86 \pm 0.49\% & \text{(BELLE [31])} \end{cases} \quad \Rightarrow \ |V_{cb}| = (41 \pm 2_{(\rm th)}) \times 10^{-3}\,. \tag{2.6}$$

Future improvements are likely to come from combined analyses using inclusive spectra to determine m_b and λ_1 (or, equivalently, $\bar\Lambda$ and λ_1). It had been suggested that moments of the $B \to X_c\ell\bar\nu$ lepton spectrum [32, 33, 34] or hadronic invariant mass spectrum [35, 34], or the $B \to X_s\gamma$ photon spectrum [36, 37] can be used to determine these parameters. Each measurement is a band in the $\bar\Lambda - \lambda_1$ plane, and the combination of several of them can pin down $\bar\Lambda$ and λ_1, and also test theoretical assumptions of the method. I.e., if quark-hadron duality were violated at the several percent level, it should show up as an inconsistency.

The first such analysis was done recently by CLEO [38, 39], using the two moments shown in Fig. 3. Combining with their semileptonic rate measurement, they obtain

$$|V_{cb}| = (40.4 \pm 1.3) \times 10^{-3}\,. \qquad (2.7)$$

The advantage of this measurement is that a sizable part of the hard-to-quantify theory error in Eq. (2.6) is traded for experimental errors on the moment measurements. To make further progress, one must quantify better the accuracy of quark-hadron duality, but if no problems are encountered $\sigma(|V_{cb}|) \sim 2\%$ may be achievable.

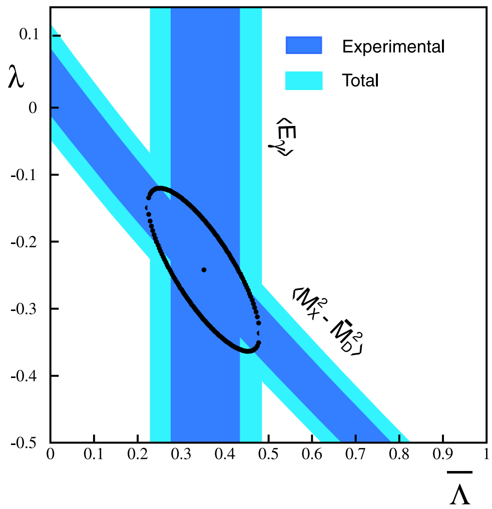

Figure 3: $\bar\Lambda$ and λ_1 from $\langle m_X^2 - \overline{m}_D^2\rangle$ in $B \to X_c\ell\bar\nu$ and $\langle E_\gamma\rangle$ in $B \to X_s\gamma$ [38].

It will continue to be important to pursue both the inclusive and exclusive measurements of $|V_{cb}|$. Since both the theoretical and the experimental systematic uncertainties are different, agreement between the two determinations will remain to be a very powerful cross-check that the errors are as well understood as claimed.

2.3 Inclusive $B \to X_u\ell\bar\nu$ spectra and $|V_{ub}|$

If it were not for the ~ 100 times larger $b \to c$ background, measuring $|V_{ub}|$ would be as "easy" as $|V_{cb}|$. The total $B \to X_u\ell\bar\nu$ rate can be predicted in the OPE with small uncertainty [28],

$$|V_{ub}| = (3.04 \pm 0.06_{(\text{pert})} \pm 0.08_{(m_b)}) \times 10^{-3} \left(\frac{\mathcal{B}(\bar B \to X_u\ell\bar\nu)}{0.001}\,\frac{1.6\,\text{ps}}{\tau_B}\right)^{1/2}, \qquad (2.8)$$

where the errors are as discussed after Eq. (2.5). The central value in Ref. [30], 3.24×10^{-3}, was later updated to 3.08×10^{-3} [40]. If this fully inclusive rate is measured without significant cuts on the phase space, then $|V_{ub}|$ can be determined with small theoretical error. It seems that measuring this rate fully inclusively may become possible using the huge data sets expected in a couple of years at the B factories.

LEP reported measurements of the inclusive rate already, giving $\mathcal{B}(b \to u\ell\bar\nu) = (1.71\pm 0.31\pm0.37\pm0.21)\times10^{-3}$ [14]. It is very hard from the outside to understand what region of the Dalitz plot these results are sensitive to. If it is the low X_u invariant mass region, then there is a sizable theoretical uncertainty (see below). As also emphasized in Refs. [41, 42], it would be most desirable to present the results also in a form which is as theory-independent as possible, and quote the rate as measured in a given kinematic region.

When kinematic cuts are used to distinguish the $b \to c$ background from the $b \to u$ signal, the behavior of the OPE can be affected dramatically. There are three qualitatively different regions of phase space, depending on the allowed invariant mass and energy (in the B rest frame) of the hadronic final state:

(i) $m_X^2 \gg E_X \Lambda_{\rm QCD} \gg \Lambda_{\rm QCD}^2$: the OPE converges, and the first few terms are expected to give reliable result. This is the case for the $B \to X_c \ell \bar\nu$ width relevant for measuring $|V_{cb}|$.

(ii) $m_X^2 \sim E_X \Lambda_{\rm QCD} \gg \Lambda_{\rm QCD}^2$: an infinite set of equally important terms in the OPE must be resummed; the OPE becomes a twist expansion and nonperturbative input is needed.

(iii) $m_X \sim \Lambda_{\rm QCD}$: the final state is dominated by resonances, and it is not known how to compute any inclusive quantity reliably.

Experimentally, there are several possibilities to remove the charm background: the charged lepton endpoint region used to first observe $b \to u$ transition, $E_\ell > (m_B^2 - m_D^2)/(2m_B)$, the low hadronic invariant mass region, $m_X < m_D$ [43, 44], and the large dilepton invariant mass region $q^2 \equiv (p_\ell + p_\nu)^2 > (m_B - m_D)^2$ [45]. These contain roughly 10%, 80%, and 20% of the rate, respectively. Measuring m_X or q^2 require reconstruction of the neutrino, which is challenging.

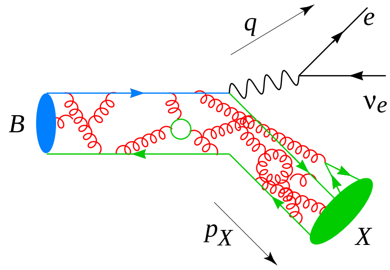

Figure 4: $B \to X \ell \bar\nu$ decay.

The problem for theory is that the phase space regions $E_\ell > (m_B^2 - m_D^2)/(2m_B)$ and $m_X < m_D$ both belong to the regime (ii), because these cuts impose $m_X \lesssim m_D$ and $E_X \lesssim m_B$, and numerically $\Lambda_{\rm QCD}\, m_B \sim m_D^2$. The region $m_X < m_D$ is better than $E_\ell > (m_B^2 - m_D^2)/(2m_B)$ inasmuch as the expected rate is a lot larger, and the inclusive description is expected to hold better. But nonperturbative input is needed, formally, at the $\mathcal{O}(1)$ level in both cases, which is why the model dependence increases rapidly if the m_X cut is lowered below m_D [43]. These regions of the Dalitz plot are shown in Fig. 5.

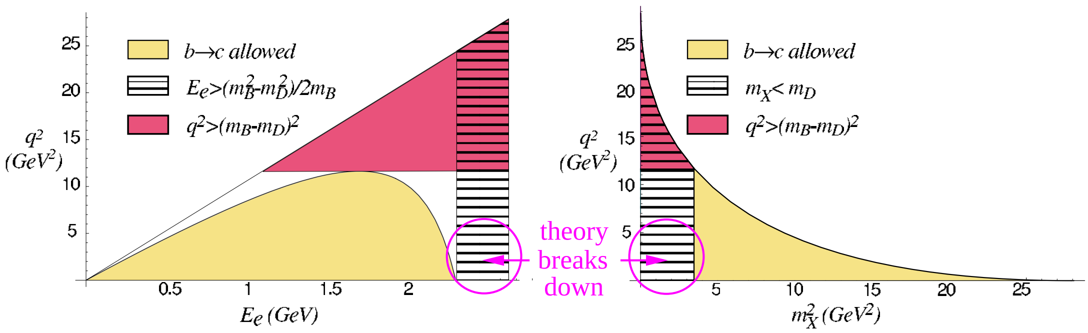

Figure 5: Dalitz plots for $B \to X \ell \bar\nu$ decay in terms of E_ℓ and q^2 (left), and m_X^2 and q^2 (right).

The nonperturbative input needed to predict the spectra in the large E_ℓ and small m_X regions, the b quark light-cone distribution function (sometimes also called shape function), is universal at leading order, and can be related to the $B \to X_s\gamma$ photon spectrum [46]. Recently these relations have been extended to the resummed next-to-leading order corrections, and applied to the large E_ℓ and small m_X regions [47]. Weighted integrals of the $B \to X_s\gamma$ photon spectrum are equal to the $B \to X_u\ell\bar\nu$ rate in the large E_ℓ or small m_X regions. There is also a sizable correction from operators other than O_7 contributing to $B \to X_s\gamma$ [48]. The dominant theoretical uncertainty in these determinations of $|V_{ub}|$ are from subleading twist contributions, which are not related to $B \to X_s\gamma$. These are suppressed by $\Lambda_{\rm QCD}/m_b$, but their size is hard to quantify, and even formulating them is nontrivial [49]. Of course, if the lepton endpoint region is found to be dominated by the π and ρ exclusive channels, then the applicability of the inclusive description may be questioned.

In contrast to the above, in the $q^2 > (m_B - m_D)^2$ region the first few terms in the OPE dominate [45]. This cut implies $E_X \lesssim m_D$ and $m_X \lesssim m_D$, and so the $m_X^2 \gg E_X\Lambda_{\rm QCD} \gg \Lambda_{\rm QCD}^2$ criterion of regime (i) is satisfied. This relies, however, on $m_c \gg \Lambda_{\rm QCD}$, and so the OPE is effectively an expansion in $\Lambda_{\rm QCD}/m_c$ [50]. The largest uncertainties come from order $\Lambda_{\rm QCD}^3/m_{c,b}^3$ nonperturbative corrections, the b quark mass, and the perturbation series. Weak annihilation (WA) suppressed by $\Lambda_{\rm QCD}^3/m_b^3$ is important, because it enters the rate as $\delta(q^2 - m_b^2)$ [51]. Its magnitude is hard to estimate, as it is proportional to the difference of two matrix elements of 4-quark operators, which vanishes in the vacuum insertion approximation. WA could be $\sim 2\%$ of the $B \to X_u\ell\bar\nu$ rate, and, in turn, $\sim 10\%$ of the rate in the $q^2 > (m_B - m_D)^2$ region. It is even more important for the lepton endpoint region, since it is also proportional to $\delta(E_\ell - m_b/2)$. Preliminary lattice results of the matrix elements suggest a smaller size [52]. Experimentally, WA can be constrained by comparing $|V_{ub}|$ measured from B^0 and $B^\pm$ decays, and by comparing the D^0 and D_s semileptonic widths [51].

Combining the q^2 and m_X cuts can significantly reduce the theoretical uncertainties [53]. The right-hand side of Fig. 5 shows that the q^2 cut can be lowered below $(m_B - m_D)^2$ by imposing an additional cut on m_X. This changes the expansion parameter from $\Lambda_{\rm QCD}/m_c$ to $m_b\Lambda_{\rm QCD}/(m_b^2 - q_{\rm cut}^2)$, resulting in a significant decrease of the uncertainties from both the perturbation series and from the nonperturbative corrections. At the same time the uncertainty from the b quark light-cone distribution function only turns on slowly. Some representative results are give in Table 2, showing that it is possible to determine $|V_{ub}|$ with a theoretical error at the $5 - 10\%$ level using up to $\sim 45\%$ of the semileptonic decays [53].

Cuts on q^2 and m_X	Fraction of events	Error of $\lvert V_{ub}\rvert$ $\delta m_b = 80/30$MeV
$6\,{\rm GeV}^2$, m_D	46%	8%/5%
$8\,{\rm GeV}^2$, $1.7\,{\rm GeV}$	33%	9%/6%
$(m_B - m_D)^2, m_D$	17%	15%/12%

Table 2: $|V_{ub}|$ from combined cuts on q^2 and m_X.

2.4 Rare B decays

Rare B decays are very sensitive probes of new physics. There are many interesting modes sensitive to different extensions of the Standard Model. For example, $B \to X_s\gamma$ provides the best bound on the charged Higgs mass in type-II two Higgs doublet model, and also constrains the parameter space of SUSY models. The photon spectrum, which is not sensitive to new physics, is important for determinations of $|V_{ub}|$ and the b quark mass, as discussed earlier. Other rare decays such as $B \to X\ell^+\ell^-$ are sensitive through the bsZ effective coupling to SUSY and left-right symmetric models. $B \to X\nu\bar{\nu}$ can probe models containing unconstrained couplings between three 3rd generation fermions [54]. In the Standard Model these decays are sensitive to CKM angles involving the top quark, complementary to $B_{d,s}$ mixing.

This last year we learned that the CKM contributions to rare decays are likely to be the dominant ones, as they probably are for CP violation in $B \to \psi K_S$. This is supported by the measurement of $\mathcal{B}(B \to X_s\gamma)$ which agrees with the SM at the 15% level [39]; the measurement of $B \to K\ell^+\ell^-$ which is in the ballpark of the SM expectation [55]; and the non-observation of direct CP violation in $b \to s\gamma$, $-0.27 < A_{CP}(B \to X_s\gamma) < 0.10$ [56] and $-0.17 < A_{CP}(B \to K^*\gamma) < 0.08$ [57] at the 90% CL, which is expected to be tiny in the SM. These results make it less likely that we will observe orders-of-magnitude enhancements of rare B decays. It is more likely that only a broad set of precision measurements will be able to find signals of new physics.

At present, inclusive rare decays are theoretically cleaner than the exclusive ones, since they are calculable in an OPE and precise multi-loop results exist. Table 3 summarizes some of the most interesting modes. The $b \to d$ rates are expected to be about a factor of $|V_{td}/V_{ts}|^2 \sim \lambda^2$ smaller than the corresponding $b \to s$ modes shown. As a guesstimate, in $b \to q\, l_1 l_2$ decays one expects $10-20\%$ K^*/ρ and $5-10\%$ K/π. A clean theoretical interpretation of the exclusive rates requires that we know the corresponding form factors. (However, CP asymmetries are independent of the form factors.) While useful relations between form factors can be derived from heavy quark symmetry, ultimately unquenched lattice calculations will be needed for clean theoretical interpretation of exclusive decays.

Decay mode	Approximate SM rate	Present status
$B \to X_s\gamma$	3.5×10^{-4}	$(3.2\pm0.5)\times10^{-4}$
$B \to X_s\nu\bar{\nu}$	4×10^{-5}	$< 7.7 \times 10^{-4}$
$B \to \tau\nu$	4×10^{-5}	$< 5.7 \times 10^{-4}$
$B \to X_s\ell^+\ell^-$	7×10^{-6}	$< 1.0 \times 10^{-5}$
$B_s \to \tau^+\tau^-$	1×10^{-6}	
$B \to X_s\tau^+\tau^-$	5×10^{-7}	
$B \to \mu\nu$	2×10^{-7}	$< 6.5 \times 10^{-6}$
$B_s \to \mu^+\mu^-$	4×10^{-9}	$< 2 \times 10^{-6}$
$B \to \mu^+\mu^-$	1×10^{-10}	$< 2.8 \times 10^{-7}$

Table 3: Some interesting rare decays.

Exclusive decays, on the other hand, are experimentally easier to measure. There have been recently some very significant theoretical developments towards understanding the relevant heavy-to-light form factors in the region of moderate q^2 (large recoil).

It was originally observed [58] that A_{FB}, the forward-backward asymmetry in $B \to K^*\ell^+\ell^-$, vanishes at a value of q^2 independent of form factor models (near $q_0^2 = 4\,\text{GeV}^2$ in the SM, see Fig. 6). This was shown to follow model independently from the large energy limit discussed below [59, 60]. One finds the following implicit equation for q_0^2

$$C_9(q_0^2) = -C_7\, \frac{2m_B m_b}{q_0^2} \left[1 + \mathcal{O}\left(\alpha_s, \frac{\Lambda_{\rm QCD}}{m_b}\right)\right]. \tag{2.9}$$

The order α_s corrections are calculable [61, 62], but one cannot reliably estimate the $\Lambda_{\rm QCD}/E_{K^*}$ terms yet. Nevertheless, these results will allow to search for new physics in A_{FB}; C_7 is known from $B \to X_s\gamma$, so the zero of A_{FB} determines C_9, which is sensitive to new physics ($C_{7,9}$ are the effective Wilson coefficients often denoted by $C_{7,9}^{\rm eff}$, and C_9 has a mild q^2-dependence).

The above simplifications occur because the 7 form factors that parameterize all $B \to$ vector meson transitions ($B \to K^*\ell^+\ell^-$, $K^*\gamma$, or $\rho\ell\bar\nu$) can be expressed in terms of only two functions, $\xi_\perp(E)$ and $\xi_\parallel(E)$, in the limit where $m_b \to \infty$ and $E_{\rho,K^*} = \mathcal{O}(m_b)$ [59]. In the same limit, the 3 form factors that parameterize decays to pseudoscalar mesons are related to one function, $\xi_P(E)$. We are just beginning to see the foundations of these ideas clarified [60], and applications worked out. E.g., the $B \to K^*\gamma$ rate can be used to constrain the $B \to \rho\ell\bar\nu$ form factors relevant for $|V_{ub}|$ [63]. The large $\mathcal{O}(\alpha_s)$ enhancement of $B \to K^*\gamma$ together with the agreement between the measured rate and the leading order prediction using light cone sum rules for the form factor imply that the form factor predictions have sizable errors or the subleading terms in $\Lambda_{\rm QCD}/E_{\rho,K^*}$ are significant [62, 64]. How well the theory can describe these processes will test some of the ingredients entering factorization in charmless B decays.

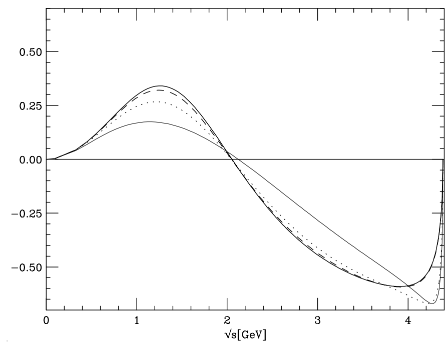

Figure 6: A_{FB} in $B \to K^*\ell^+\ell^-$ in different form factor models ($s \equiv q^2$) [58].

2.5 Semileptonic and rare decays — Summary

- $|V_{cb}|$ is known at the $\lesssim 5\%$ level; error may become half of this in the next few years using both inclusive and exclusive measurements. The inclusive requires precise determination of m_b using various spectra and tests of duality, the exclusive will rely on the lattice.

- Situation for $|V_{ub}|$ may become similar to present $|V_{cb}|$. For a precise inclusive measurement the neutrino reconstruction to obtain q^2 and m_X seems crucial (and determining m_b as mentioned above); the exclusive will require unquenched lattice calculations.

- Important progress towards understanding exclusive rare decays in the small q^2 regime, $B \to \rho\ell\bar\nu$, $K^{(*)}\gamma$, and $K^{(*)}\ell^+\ell^-$ below the ψ. This increases the sensitivity to new physics, and may also test some ingredients entering factorization in charmless decays.

3. Nonleptonic decays

In this Section I discuss three topics where important developments occurred recently. The first is factorization in exclusive hadronic B decays. Especially charmless decays are very important for studying CP violation. The second is inclusive widths and lifetimes, where OPE based calculations are possible. The third is $D-\overline{D}^0$ mixing, where there have also been new experimental and theoretical results.

3.1 Factorization in exclusive B decays

Until recently very little was known model independently about exclusive nonleptonic B decays. Crudely speaking, factorization is the hypothesis that, starting from the effective nonleptonic Hamiltonian, one can estimate matrix elements of four-quark operators by grouping the quark fields into a pair that can mediate $B \to M_1$ decay (M_1 inherits the "brown muck" of the decaying B), and another pair that can describe vacuum $\to M_2$ transition. E.g., in $\overline{B}^0 \to D^{(*)+}\pi^-$, this amounts to the assumption that the contributions of gluons "parallel" to the W are calculable perturbatively or suppressed by $\Lambda_{\rm QCD}/m_Q$ (see Fig. 7).

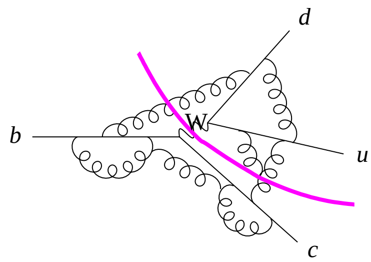

Figure 7: Sketch of factorization in $\overline{B}^0 \to D^{(*)+}\pi^-$ decay.

It has long been known that if M_1 is heavy and M_2 is light, such as $\overline{B}^0 \to D^{(*)+}\pi^-$, then "color transparency" may justify factorization [65, 66, 67]. The physical picture is that the two quarks forming the π must emerge from the short distance process in a small color dipole state (two fast collinear quarks in a color singlet), and at the same time the wave function of the brown muck in the B only changes moderately since the D recoil is small. Recently it was shown to 2-loops [68], and to all orders in perturbation theory [69], that in such decays factorization is the leading result in a systematic expansion in powers of $\alpha_s(m_Q)$ and $\Lambda_{\rm QCD}/m_Q$. While the α_s corrections are calculable, little is known from first principles about those suppressed, presumably, by powers of $\Lambda_{\rm QCD}/m_b$. A renormalon analysis suggests that in $\overline{B}^0 \to D^{(*)+}\pi^-$, where the light-cone wave function of M_2 (the π) is symmetric, nonperturbative corrections are actually suppressed by two powers [70].

It is important to test experimentally how well factorization works, and learn about the size of power suppressed effects. The $\langle D^{(*)}|\bar{c}_L\gamma^\mu b_L|\overline{B}^0\rangle$ matrix element is measured in semileptonic $B \to D^{(*)}$ decay, while $\langle X|\bar{u}_L\gamma^\mu d_L|0\rangle$ for $X=\pi,\rho$ is given by the known decay constants. Thus, in "color allowed" decays, such as $\overline{B}^0 \to D^{(*)+}\pi^-$ and $D^{(*)+}\rho^-$, factorization has been observed to work at the $\sim 10\%$ level. These tests get really interesting just around this level, since there is another argument that supports factorization, which is independent of the heavy mass limit. It is the large N_c limit ($N_c = 3$ is the number of colors), which implies for such decays that factorization violation is suppressed by $1/N_c^2$. The large N_c argument for factorization is independent of the final state, whereas the one based on the heavy quark limit predicts that the accuracy depends on the kinematics of the decay.

One of the predictions of QCD factorization in $B \to D\pi$ is that amplitudes involving the spectator quark in the B going into the π should be power suppressed [68], and therefore,

$$\mathcal{B}(B \to D^{(*)0}\pi^-)/\mathcal{B}(B \to D^{(*)+}\pi^-) = 1 + \mathcal{O}(\Lambda_{\rm QCD}/m_Q)\,. \tag{3.1}$$

However, experimentally, this ratio is in the ballpark of 1.8 with errors around 0.3 for both D and D^* and also for π replaced by ρ. This has been argued to be due to $\mathcal{O}(\Lambda_{\rm QCD}/m_c)$ corrections, which may be sizable [68].

$\mathcal{B}(B \to D^0\pi^0)$	$\mathcal{B}(B \to D^{*0}\pi^0)$	$[\times 10^{-4}]$
$3.1 \pm 0.4 \pm 0.5$	$2.7^{+0.8}_{-0.7}{}^{+0.5}_{-0.6}$	BELLE [71]
$2.74^{+0.36}_{-0.32} \pm 0.55$	$2.20^{+0.59}_{-0.52} \pm 0.79$	CLEO [72]

Table 4: Color suppressed $B \to D^{(*)0}\pi^0$ branching ratios.

The first observations of "color suppressed" B decays, $B \to D^{(*)0}\pi^0$, were reported at this conference. The results are summarized in Table 4. These rates are larger than earlier theoretical expectations (or than the upper bound for $D^0\pi^0$ in the Y2K PDG). This data allows, for the first time, to extract the strong phase difference between the $\Delta I = \frac{3}{2}$ and $\frac{1}{2}$ amplitudes from the measured $B \to D^+\pi^-$, $D^0\pi^-$, and $D^0\pi^0$ rates. Factorization predicts that this phase should be power suppressed. My slides at the conference showed that this phase was around 24° with asymmetric errors around 6°. Since then, several analyses are published with varying conclusions about the meaning of these results [73]. It will be interesting to see what happens when the experimental errors decrease.

There are many other testable predictions. E.g., factorization also holds in $\overline{B}{}^0 \to D^{(*)+}D_s^{(*)-}$ within the (presently sizable) errors, which is interesting because the heavy D_s meson must come from the W boson [74]. At some level one expects to see deviations from factorization in this decay which are larger than those in $\overline{B}{}^0 \to D^{(*)+}\pi^-$. When the $B \to \pi$ semileptonic form factors and the $\overline{B}{}^0 \to \pi^+D_s^{(*)-}$ rate will be measured, it will be interesting to compare the accuracy of factorization with that in $\overline{B}{}^0 \to D^{(*)+}\pi^-$. In QCD factorization $\overline{B}{}^0 \to \pi^+D_s^{(*)-}$ is power suppressed, so corrections to "naive factorization" are not subleading in the power counting. So I would not trust $|V_{ub}|$ extracted from this rate.

It was also observed that in $B \to D^{(*)}X$, where X is a meson with spin greater than one or has a small decay constant (such as the a_0, b_1, etc., which can only be created by the weak current due to isospin breaking), the leading factorizable term vanishes, but there is a calculable $\mathcal{O}(\alpha_s)$ contribution [75]. Unfortunately there are also power suppressed uncalculable corrections, which may be comparable. Such ideas could also be useful for CP violation studies in charmless decays, by suppressing certain tree amplitudes [75, 76]. There may be preliminary evidence for one such decay, $B \to a_0\pi$ [77].

Multi-body $B \to D^{(*)}X$ modes have also been used to study corrections to factorization [78]. The advantage compared to two-body channels is that the accuracy of factorization can be studied for a final state with fixed particle content, by examining the differential decay rate as a function of the invariant mass of the light hadronic state X (this was also suggested in Refs. [66, 79]). If factorization works primarily due to the large N_c limit then its accuracy is not expected to decrease as the X invariant mass, m_X, increases. If factorization is mostly due to perturbative QCD then there should be corrections which

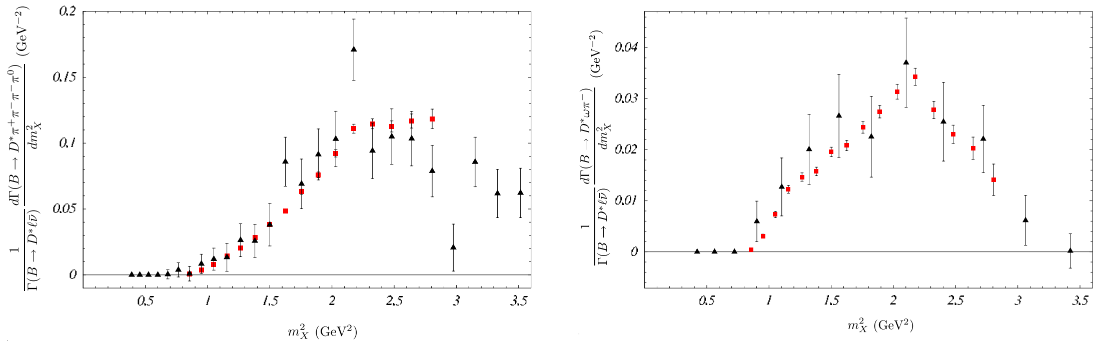

Figure 8: $\mathrm{d}\Gamma(B \to D^*\pi^+\pi^-\pi^-\pi^0)/\mathrm{d}m_X^2$, where m_X is the $\pi^+\pi^-\pi^-\pi^0$ invariant mass (left), and $\mathrm{d}\Gamma(B \to D^*\omega\pi)/\mathrm{d}m_X^2$, where m_X is the $\omega\pi^-$ invariant mass (right), normalized to the $B \to D^*\ell\bar\nu$ rate. Black triangles are B decay data, red squares are the predictions using τ data [78].

grow with m_X. Combining data for hadronic τ decays and semileptonic B decays allows such tests to be made for a variety of final states. Fig. 8 shows the comparison of the $B \to D^*\pi^+\pi^-\pi^-\pi^0$ and $D^*\omega\pi^-$ data [80] with the τ decay data [81]. The kinematic range accessible in $\tau \to 4\pi$ corresponds to $0.4 \lesssim m_{4\pi}/E_{4\pi} \lesssim 0.7$ in $B \to 4\pi$ decay. A background to these comparisons is that one or more of the pions may arise from the $\bar c_L\gamma^\mu b_L$ current creating a nonresonant $D^* + n\,\pi$ $(1 \le n \le 3)$ state or a higher D^{**} resonance. In the $\omega\pi^-$ mode this is very unlikely to be significant [78]. In the $\pi^+\pi^-\pi^-\pi^0$ mode such backgrounds can be constrained by measuring $B \to D^*\pi^+\pi^+\pi^-\pi^-$, since $\pi^+\pi^+\pi^-\pi^-$ cannot come from the $\bar u_L\gamma^\mu d_L$ current. CLEO found $\mathcal{B}(B \to D^*\pi^+\pi^+\pi^-\pi^-)/\mathcal{B}(B \to D^*\pi^+\pi^-\pi^-\pi^0) < 0.13$ at 90% CL in the $m_X^2 < 2.9\,\mathrm{GeV}^2$ region [82]. With more precise data, observing deviations that grow with m_X would be evidence that perturbative QCD is an important part of the success of factorization in $B \to D^*X$.

Calculating B decay amplitudes to charmless two-body final states is especially important for the study of CP violation. There are two approaches to these decays. BBNS [83] assume that Sudakov suppression is not effective at the B mass scale in the endpoint regions of quark distribution functions, while Keum *et al.* [84] assume that it is. They yield different power counting and often different phenomenological predictions. In the former approach the $B \to \pi\ell\bar\nu$ form factors are nonperturbative functions to be determined from data, while they are calculable in the latter. My guess would be that they are not calculable (they would be if m_b were huge), but it will take time to really decide this using data. Predictions for direct CP violation are often smaller in the former than in the latter approach. An outstanding open theoretical question is the complete formulation of power suppressed corrections. Some of them are known to be large, e.g., the "chirally enhanced" terms proportional to $m_K^2/(m_s m_b)$ which are not enhanced by any parameter of QCD in the chiral limit, just "happen to be" large, and the uncertainty related to controlling the infrared sensitivity in annihilation contributions. (See also the discussion of these issues in Refs. [41, 42].) It has also been claimed that the effects of charm loops are bigger than given by perturbation theory [85].

3.2 Inclusive nonleptonic decays, b hadron lifetimes

Inclusive nonleptonic decay rates of heavy hadrons can also be computed in an OPE, like inclusive semileptonic rates. The crucial difference is that the OPE has to be performed in the physical region, and so lifetime predictions rely on local duality, whereas inclusive semileptonic rates only rely on global duality. Formally, they are expected to have similar accuracy in the $m_b \to \infty$ limit, but it is quite possible that the scale at which local duality becomes a good approximation is larger than that for global duality. It would not be surprising if the predictions of the OPE work better for semileptonic than for nonleptonic rates.

The most recent world average b hadron lifetime ratios, together with theoretical expectations, are shown in Fig. 9 [86]. The lifetime differences are expected to be dominated by matrix elements of four-quark operators at order $(\Lambda_{\rm QCD}/m_b)^3$, which have to be determined from lattice QCD [52, 87]. For now, the smallness of $\tau(\Lambda_b)$ remains hard to explain. However, this is not an indication that semileptonic widths have similar theoretical uncertainties. The semileptonic width are in fact consistent, $\mathcal{B}(\Lambda_b \to X\ell\bar\nu)/\tau(\Lambda_b) \simeq \mathcal{B}(B \to X\ell\bar\nu)/\tau(B)$, within the $\sim 15\%$ error of the present experimental data [14].

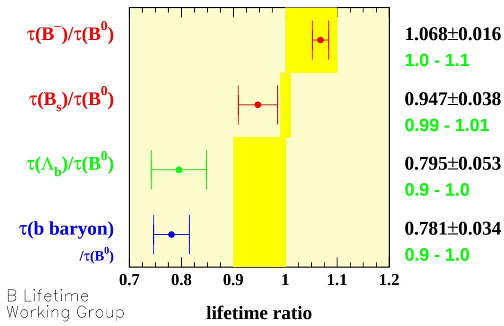

Figure 9: b hadron lifetime ratios [86].

The assumption in the OPE calculation of nonleptonic widths related to local duality has been questioned recently. In the 't Hooft Model (two dimensional QCD) it was found numerically that the widths of a heavy meson and a heavy quark differ by order $\Lambda_{\rm QCD}/m_Q$, and one needs to do an (unphysical) smearing over m_Q to reduce the discrepancy to $\Lambda_{\rm QCD}^2/m_Q^2$ [88].

3.3 $D^0 - \overline{D}{}^0$ mixing

The D^0 system is unique among the neutral mesons in that it is the only one whose mixing proceeds via intermediate states with down-type quarks. $D^0 - \overline{D}{}^0$ mixing is a sensitive probe of new physics, because the SM prediction for $x \equiv \Delta M_D/\Gamma_D$, $y \equiv \Delta\Gamma_D/2\Gamma_D$, and the CP violating phase in the mixing, ϕ, are very small. While y is expected to be dominated by SM processes, x and ϕ could be significantly enhanced by new physics.

D^0 mixing is very slow in the SM, because the third generation plays a negligible role due to the smallness of $|V_{ub}V_{cb}|$, the GIM cancellation is very effective due to the smallness of m_b/m_W, and the mixing is also suppressed by $SU(3)$ breaking. x and y are hard to estimate reliably because the charm quark is neither heavy

ratio	4-quark	6-quark	8-quark
$\dfrac{\Delta M}{\Delta M_{\rm box}}$	1	$\dfrac{\Lambda^2}{m_s m_c}$	$\dfrac{\alpha_s}{4\pi}\,\dfrac{\Lambda^4}{m_s^2 m_c^2}$
$\dfrac{\Delta\Gamma}{\Delta M}$	$\dfrac{m_s^2}{m_c^2}$	$\dfrac{\alpha_s}{4\pi}$	$\dfrac{\alpha_s}{4\pi}\,\beta_0$

Table 5: ΔM and $\Delta\Gamma$ in the OPE ($\Lambda \sim 1\,$GeV).

enough to trust the "inclusive" approach based on the OPE, nor light enough to trust the "exclusive" approach which sums over intermediate hadronic states. The short distance box diagram contributes $x_{\rm box} \sim$ few $\times 10^{-5}$ since it is suppressed by $m_s^4/(m_W^2 m_c^2)$, and $y_{\rm box} \sim$ few $\times 10^{-7}$ since it has an additional m_s^2/m_c^2 helicity suppression. Higher order terms in the OPE are very important, because they are suppressed by fewer powers of m_s (see Table 5) [89, 90, 91]. With large uncertainties due to the hadronic matrix elements, most estimates yield $x, y \lesssim 10^{-3}$.

There are three types of experiments which measure x and y. Each is actually sensitive to a combination of x and y, rather than to either quantity directly. First, the D^0 lifetime difference to CP even and CP odd final states can be measured by comparing the lifetimes to a flavor and a CP eigenstate. To leading order,

Value of y_{CP}	Experiment
$0.8 \pm 3.1\%$	E791 [92]
$3.4 \pm 1.6\%$	FOCUS [93]
$-1.1 \pm 2.9\%$	CLEO [94]
$-0.5 \pm 1.3\%$	BELLE [95]
$-1.0 \pm 2.8\%$	BABAR [96]

Table 6: y_{CP} measurements.

$$y_{CP} = \frac{\tau(D \to \pi^+ K^-)}{\tau(D \to K^+ K^-)} - 1 = y \cos\phi - x \sin\phi \, \frac{A_m}{2} \,, \tag{3.2}$$

where $A_m = |q/p|^2 - 1$, which is very small in the SM. The present data in Table 6 yield a world average $y_{CP} \simeq 0.65 \pm 0.85\%$. Second, the time dependence of doubly Cabibbo suppressed decays, such as $D^0 \to K^+\pi^-$ [97], is sensitive to the three quantities

$$(x\cos\delta + y\sin\delta)\cos\phi \,, \qquad (y\cos\delta - x\sin\delta)\sin\phi \,, \qquad x^2 + y^2 \,, \tag{3.3}$$

where δ is the strong phase between the Cabibbo allowed and doubly Cabibbo suppressed amplitudes (see Fig. 10). A similar study for $D^0 \to K^-\pi^+\pi^0$ would be valuable, with the strong phase difference extracted simultaneously from the Dalitz plot analysis [98]. Third, one can search for D mixing in semileptonic decays [99], which is sensitive to $x^2 + y^2$.

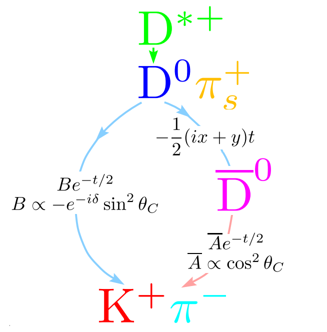

Figure 10: $D^0 \to K^+\pi^-$.

Although y is expected to be determined by Standard Model processes, its value affects significantly the sensitivity to new physics [100]. If y is larger or much larger than x, then the observable CP violation in D^0 mixing is necessarily small, even if new physics dominates x. A recent estimate of y calculated $SU(3)$ breaking in phase space differences, and found that $y \sim 1\%$ can easily be accommodated in the SM [91]. Final states in D^0 decay can be decomposed in representations of $SU(3)$. The cancellation between decays to members of a given representation can be significantly violated because the final states containing larger number of strange hadrons have smaller phase space, or can even be completely forbidden. Such effects might enhance y more significantly than they affect x.

Therefore, searching for new physics and CP violation in $D^0 - \overline{D}{}^0$ mixing should aim at precise measurements of both x and y, and at more complicated analyses involving the extraction of the strong phase in the time dependence of doubly Cabibbo suppressed decays.

3.4 Nonleptonic decays — Summary

- In nonleptonic $B \to D^{(*)}X$ decay, where X is a low mass hadronic state, factorization has been established in the heavy quark limit, at leading order in $\Lambda_{\rm QCD}/m_Q$.

- Flood of new and more precise data will allow many tests of factorization and tell us the significance of unknown power suppressed terms, hopefully also in charmless decays.

- In the D system the only unambiguous signal of new physics is CP violation; observation of a large Δm_D can only be a clear sign if $\Delta\Gamma_D$ is smaller, so crucial to measure both.

4. Conclusions

I was not supposed to talk about CP violation, but I had no chance to succeed, because in order to test the Standard Model in flavor physics all possible clean measurements which give model independent information on short distance parameters are very important, whether CP violating or conserving.

With the recent fairly precise measurement of $\sin 2\beta$ and other data, the CKM contributions to flavor physics and CP violation are likely to be the dominant ones. The next goal is not simply to measure ρ and η, or α and γ, but to probe the flavor sector of the SM until it breaks. This can be hoped to be achieved in B decays by overconstraining measurements of the unitarity triangle. Measurements which are redundant in the SM but sensitive to different short distance physics are also very important, since correlations may give information on the new physics we are encountering (e.g., comparing $\Delta m_s/\Delta m_d$ with $\mathcal{B}(B \to X_s\ell^+\ell^-)/\mathcal{B}(B \to X_d\ell^+\ell^-)$ is not "just another way" to measure $|V_{ts}/V_{td}|$).

In many cases hadronic uncertainties are significant and hard to quantify. The sensitivity to new physics and the accuracy with which the SM can be tested will depend on our ability to disentangle the short distance physics from nonperturbative effects of hadronization. While we all want smaller errors, ϵ'_K reminds us to be conservative with theoretical uncertainties. One theoretically clean measurement is worth ten dirty ones. But it does change with time what is theoretically clean, and I hope to have conveyed that there are significant recent developments towards understanding the hadronic physics crucial both for standard model measurements and for searches for new physics. For example, (i) for the determination of $|V_{ub}|$ from inclusive B decay; (ii) for understanding exclusive rare decay form factors at small q^2; and (iii) for establishing factorization in certain nonleptonic decays.

In testing the SM and searching for new physics, our understanding of CKM parameters and hadronic physics will have to improve in parallel, since except for a few clean cases (like $\sin 2\beta$) the theoretical uncertainties can be reduced by doing several measurements, or by learning from comparisons with data how accurate certain theoretical assumptions are. In some cases data will help to constrain or get rid of nasty things hard to know model independently from theory (e.g., excited state contributions to certain processes).

With the recent spectacular start of the B factories an exciting era in flavor physics has begun. The precise measurements of $\sin 2\beta$ together with the sides of the unitarity triangle, $|V_{ub}/V_{cb}|$ at the e^+e^- B factories and $|V_{td}/V_{ts}|$ at the Tevatron, will allow to observe small deviations from the Standard Model. The large statistics will allow the study of rare decays and to improve sensitivity to observables which vanish in the SM (e.g., certain CP asymmetries); these measurements have individually the potential to discover physics beyond the SM. If new physics is seen, then a broad set of measurements at both e^+e^- and hadronic B factories and $K \to \pi\nu\bar{\nu}$ may allow to discriminate between various scenarios. This is a vibrant theoretical and experimental program, and I think the most concise summary of the status of the field is:

> "This is not the end. It is not even the beginning of the end.
> But it is, perhaps, the end of the beginning."
> W. Churchill (Nov. 10, 1942)

Acknowledgments

It is a pleasure to acknowledge numerous discussions with Adam Falk, Mike Luke, Helen Quinn, and Mark Wise, and many colleagues for help with the recent data, especially Karl Ecklund, Gautier Hamel de Monchenault, and Vivek Sharma. I would like to thank the organizers for the invitation to present this talk and for their support. This work was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, Division of High Energy Physics, of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

References

[1] G. Hamel de Monchenault, plenary talk, these Proceedings, hep-ex/0112007.

[2] C. Weiser, these Proceedings; J. Thom, these Proceedings.

[3] A.B. Wicklund, these Proceedings.

[4] S. Ryan, hep-lat/0111010; and references therein.

[5] B. Grinstein *et al.*, Nucl. Phys. B380 (1992) 369.

[6] N. Yamada *et al.*, JLQCD Collaboration, hep-lat/0110087.

[7] M. Beneke, plenary talk, these Proceedings.

[8] A. Hocker, H. Lacker, S. Laplace and F. Le Diberder, Eur. Phys. J. C21 (2001) 225; and updates at `http://www.slac.stanford.edu/~laplace/ckmfitter.html`.

[9] N. Isgur and M.B. Wise, Phys. Lett. B232 (1989) 113; Phys. Lett. B237 (1990) 527.

[10] A. Czarnecki, Phys. Rev. Lett. 76 (1996) 4124;
A. Czarnecki and K. Melnikov, Nucl. Phys. B505 (1997) 65.

[11] M.E. Luke, Phys. Lett. B252 (1990) 447.

[12] M. Neubert, Z. Ligeti and Y. Nir, Phys. Lett. B301 (1993) 101; Phys. Rev. D47 (1993) 5060;
Z. Ligeti, Y. Nir and M. Neubert, Phys. Rev. D49 (1994) 1302.

[13] S. Hashimoto *et al.*, Phys. Rev. D61 (2000) 014502;

[14] F. Palla, these Proceedings, and references therein.

[15] J.P. Alexander *et al.*, CLEO Collaboration, hep-ex/0007052.

[16] K. Abe *et al.*, BELLE Collaboration, hep-ex/0111060.

[17] S. Hashimoto *et al.*, hep-ph/0110253.

[18] C.G. Boyd, B. Grinstein, R.F. Lebed, Phys. Lett. B353 (1995) 306; Nucl. Phys. B461 (1996) 493; Phys. Rev. D56 (1997) 6895.

[19] I. Caprini, L. Lellouch and M. Neubert, Nucl. Phys. B530 (1998) 153.

[20] A. Le Yaouanc *et al.*, Phys. Lett. B520 (2001) 25; Phys. Lett. B520 (2001) 59.

[21] N. Uraltsev, Phys. Lett. B501 (2001) 86.

[22] A.F. Falk, Nucl. Phys. B378 (1992) 79.

[23] A.K. Leibovich, Z. Ligeti, I.W. Stewart and M.B. Wise, Phys. Rev. Lett. 78 (1997) 3995; Phys. Rev. D57 (1998) 308.

[24] J. Bartelt *et al.*, CLEO Collaboration, Phys. Rev. Lett. 82 (1999) 3746.

[25] K. Abe *et al.*, BELLE Collaboration, hep-ex/0111082.

[26] B. Grinstein and Z. Ligeti, hep-ph/0111392.

[27] J. Chay, H. Georgi and B. Grinstein, Phys. Lett. B247 (1990) 399;
M.A. Shifman and M.B. Voloshin, Sov. J. Nucl. Phys. 41 (1985) 120;
I.I. Bigi, N.G. Uraltsev and A.I. Vainshtein, Phys. Lett. B293 (1992) 430 [E. B297 (1992) 477];
I.I. Bigi, M.A. Shifman, N.G. Uraltsev and A.I. Vainshtein, Phys. Rev. Lett. 71 (1993) 496;
A.V. Manohar and M.B. Wise, Phys. Rev. D49 (1994) 1310.

[28] A.H. Hoang, Z. Ligeti and A.V. Manohar, Phys. Rev. Lett. 82 (1999) 277; Phys. Rev. D59 (1999) 074017.

[29] A.H. Hoang, Phys. Rev. D61 (2000) 034005; M. Beneke and A. Signer, Phys. Lett. B471 (1999) 233; K. Melnikov and A. Yelkhovsky, Phys. Rev. D59 (1999) 114009.

[30] I.I. Bigi, M.A. Shifman and N. Uraltsev, Ann. Rev. Nucl. Part. Sci. 47 (1997) 591.

[31] E. Won, these Proceedings, BELLE-CONF-0123.

[32] M.B. Voloshin, Phys. Rev. D51 (1995) 4934.

[33] M. Gremm, A. Kapustin, Z. Ligeti and M.B. Wise, Phys. Rev. Lett. 77 (1996) 20;
M. Gremm and I. Stewart, Phys. Rev. D55 (1997) 1226.

[34] M. Gremm and A. Kapustin, Phys. Rev. D55 (1997) 6924.

[35] A.F. Falk, M. Luke, and M.J. Savage, Phys. Rev. D53 (1996) 2491; Phys. Rev. D53 (1996) 6316.

[36] A. Kapustin and Z. Ligeti, Phys. Lett. B355 (1995) 318.

[37] Z. Ligeti, M. Luke, A.V. Manohar and M.B. Wise, Phys. Rev. D60 (1999) 034019.

[38] D. Cronin-Hennessy *et al.*, CLEO Collaboration, hep-ex/0108033.

[39] S. Chen *et al.*, CLEO Collaboration, hep-ex/0108032;

[40] N. Uraltsev, Int. J. Mod. Phys. A14 (1999) 4641.

[41] H. Quinn, hep-ph/0111169.

[42] M.B. Wise, hep-ph/0111167.

[43] A.F. Falk, Z. Ligeti and M.B. Wise, Phys. Lett. B406 (1997) 225.

[44] I. Bigi, R.D. Dikeman and N. Uraltsev, Eur. Phys. J. C4 (1998) 453.

[45] C.W. Bauer, Z. Ligeti and M. Luke, Phys. Lett. B479 (2000) 395; see also: hep-ph/0007054.

[46] M. Neubert, Phys. Rev. D49 (1994) 4623;
I.I. Bigi, M.A. Shifman, N.G. Uraltsev and A.I. Vainshtein, Int. J. Mod. Phys. A9 (1994) 2467.

[47] A.K. Leibovich and I.Z. Rothstein, Phys. Rev. D61 (2000) 074006;
A.K. Leibovich, I. Low and I.Z. Rothstein, Phys. Rev. D61 (2000) 053006; Phys. Lett. B486 (2000) 86; Phys. Lett. B513 (2001) 83.

[48] M. Neubert, Phys. Lett. B513 (2001) 88.

[49] C.W. Bauer, M. Luke and T. Mannel, hep-ph/0102089.

[50] M. Neubert, JHEP 0007 (2000) 022; M. Neubert and T. Becher, hep-ph/0105217.

[51] M.B. Voloshin, Phys. Lett. B515 (2001) 74.

[52] M. Di Pierro and C.T. Sachrajda, Nucl. Phys. B534 (1998) 373;

[53] C.W. Bauer, Z. Ligeti and M. Luke, Phys. Rev. D64 (2001) 113004.

[54] Y. Grossman, Z. Ligeti and E. Nardi, Nucl. Phys. B465 (1996) 369 [E. B480 (1996) 753].

[55] K. Abe *et al.*, BELLE Collaboration, hep-ex/0109026.

[56] T.E. Coan *et al.*, CLEO Collaboration, Phys. Rev. Lett. 86 (2001) 5661.

[57] B. Aubert *et al.*, BABAR Collaboration, hep-ex/0110065.

[58] G. Burdman, Phys. Rev. D57 (1998) 4254.

[59] J. Charles *et al.*, Phys. Rev. D60 (1999) 014001.

[60] C.W. Bauer, S. Fleming and M.E. Luke, Phys. Rev. D63 (2001) 014006;
C.W. Bauer, S. Fleming, D. Pirjol and I.W. Stewart, Phys. Rev. D63 (2001) 114020.

[61] M. Beneke and T. Feldmann, Nucl. Phys. B592 (2001) 3.

[62] M. Beneke, T. Feldmann and D. Seidel, Nucl. Phys. B612 (2001) 25.

[63] G. Burdman and G. Hiller, Phys. Rev. D63 (2001) 113008.

[64] S.W. Bosch and G. Buchalla, hep-ph/0106081.

[65] J.D. Bjorken, Nucl. Phys. Proc. Suppl. 11 (1989) 325;

[66] M.J. Dugan and B. Grinstein, Phys. Lett. B255 (1991) 583.

[67] H.D. Politzer and M.B. Wise, Phys. Lett. B257 (1991) 399.

[68] M. Beneke, G. Buchalla, M. Neubert and C.T. Sachrajda, Nucl. Phys. B591 (2000) 313.

[69] C.W. Bauer, D. Pirjol and I.W. Stewart, Phys. Rev. Lett. 87 (2001) 201806; hep-ph/0109045.

[70] C.N. Burrell and A.R. Williamson, Phys. Rev. D64 (2001) 034009; T. Becher, M. Neubert and B.D. Pecjak, hep-ph/0102219.

[71] K. Abe *et al.*, BELLE Collaboration, hep-ex/0109021.

[72] T.E. Coan *et al.*, CLEO Collaboration, hep-ex/0110055.

[73] Z–z. Xing, hep-ph/0107257; H.Y. Cheng, hep-ph/0108096; M. Neubert and A.A. Petrov, Phys. Lett. B519 (2001) 50; J.P. Lee, hep-ph/0109101.

[74] Z. Luo and J.L. Rosner, Phys. Rev. D64 (2001) 094001.

[75] M. Diehl and G. Hiller, JHEP 0106 (2001) 067; Phys. Lett. B517 (2001) 125.

[76] S. Laplace and V. Shelkov, hep-ph/0105252.

[77] B. Aubert *et al.*, BABAR Collaboration, hep-ex/0107075.

[78] Z. Ligeti, M. Luke and M.B. Wise, Phys. Lett. B507 (2001) 142.

[79] C. Reader and N. Isgur, Phys. Rev. D47 (1993) 1007.

[80] J.P. Alexander *et al.*, CLEO Collaboration, Phys. Rev. D64 (2001) 092001.

[81] K.W. Edwards *et al.*, CLEO Collaboration, Phys. Rev. D61 (2000) 072003.

[82] K.W. Edwards *et al.*, CLEO Collaboration, hep-ex/0105071.

[83] M. Beneke, G. Buchalla, M. Neubert and C.T. Sachrajda, Phys. Rev. Lett. 83 (1999) 1914; Nucl. Phys. B606 (2001) 245.

[84] Y.Y. Keum, H–n. Li and A.I. Sanda, Phys. Lett. B504 (2001) 6; Phys. Rev. D63 (2001) 054008; Y.Y. Keum and H–n. Li, Phys. Rev. D63 (2001) 074006.

[85] M. Ciuchini *et al.*, Phys. Lett. B515 (2001) 33; S.J. Brodsky and S. Gardner, hep-ph/0108121.

[86] K. Osterberg, these Proceedings.

[87] M. Di Pierro, C.T. Sachrajda and C. Michael, Phys. Lett. B468 (1999) 143; D. Becirevic, hep-ph/0110124.

[88] B. Grinstein, Phys. Rev. D64 (2001) 094004; and references therein.

[89] H. Georgi, Phys. Lett. B297 (1992) 353.

[90] I. Bigi and N. Uraltsev, Nucl. Phys. B592 (2000) 92.

[91] A.F. Falk, Y. Grossman, Z. Ligeti and A.A. Petrov, hep-ph/0110317.

[92] E.M. Aitala *et al.*, E791 Collaboration, Phys. Rev. Lett. 83 (1999) 32.

[93] J.M. Link *et al.*, FOCUS Collaboration, Phys. Lett. B485 (2000) 62.

[94] D. Cronin-Hennessy *et al.*, CLEO Collaboration, hep-ex/0102006.

[95] K. Abe *et al.*, BELLE Collaboration, hep-ex/0111026.

[96] B. Aubert *et al.*, BABAR Collaboration, hep-ex/0109008.

[97] R. Godang *et al.*, CLEO Collaboration, Phys. Rev. Lett. 84 (2000) 5038.

[98] G. Brandenburg *et al.*, CLEO Collaboration, Phys. Rev. Lett. 87 (2001) 071802 .

[99] E.M. Aitala *et al.*, E791 Collaboration, Phys. Rev. D57 (1998) 13.

[100] S. Bergmann *et al.*, Phys. Lett. B486 (2000) 418.