I discuss two aspects of top-quark physics: $V_{tb}$ and the top-quark mass. The similarities and differences with bottom and charm physics are emphasized.

While preparing this talk, I was struck by how different the physics of the top quark is from that of bottom and charm. This is largely a consequence of the very short lifetime of the top quark, $\Gamma^{-1} \approx (1.5 \text{ GeV})^{-1}$, less than the characteristic time scale of nonperturbative QCD, $\Lambda_{QCD}^{-1} \approx (200 \text{ MeV})^{-1}$. As a result, the physics of the top quark is free of the effects of nonperturbative QCD. Thus one is not concerned with form factors, decay constants, exclusive decays, and other such topics which one usually associates with heavy-flavor physics. The lack of nonperturbative effects allows for precision studies of the top quark, which gives us access to interesting physics both within the standard model and, hopefully, beyond.

Within the standard model, there are only a few parameters associated with the top quark: $m_t, V_{tb}, V_{ts}$, and $V_{td}$. One goal of top-quark physics is to measure these parameters precisely. However, we have much higher hopes for top-quark physics; we hope that it will reveal physics beyond the standard model.

In this talk I will consider top physics within the standard model. If we are to find physics beyond the standard model, it is essential that we first understand top physics within the standard model. Furthermore, I have chosen to concentrate on two topics which I believe to be of particular interest to heavy-flavor physicists. These topics are $V_{tb}$, and the measurement of $m_t$ from the $t\bar{t}$ threshold in lepton colliders.

1 $V_{tb}$

It is amusing that, within the context of three generations, $V_{tb}$ is the best-known CKM matrix element (as a percentage of its value), $V_{tb} = 0.9989 - 0.9993$. Despite the fact that it has never been directly measured. We owe this to beauty physicists, who have measured $V_{cb}$ and $V_{ub}$ to be small: three-generation

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\textsuperscript{a} Assuming three generations, $V_{ts}$ and $V_{td}$ are determined indirectly from loop processes in $B$ and $K$ physics. 

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unitarity then implies that $V_{tb}$ must be near unity. The direct measurement of $V_{tb}$ is therefore only of interest if we entertain the possibility of more than three generations. In this case $V_{tb}$ is almost entirely unconstrained, $V_{tb} = 0 - 0.9993.$

We heard in the previous talk that CDF has measured $V_{tb} > 0.76$ (95% C.L.), assuming three generations. It is important to realize that this bound disappears if there are more than three generations. Recall that CDF measures the ratio of the branching ratio of the top quark to bottom quarks over the branching ratio to all quarks,

$$\frac{BR(t \to bW)}{BR(t \to qW)} = \frac{|V_{tb}|^2}{|V_{td}|^2 + |V_{ts}|^2 + |V_{tb}|^2} = .99 \pm .29.$$  \hfill (1)

If there are just three generations, then the denominator in the second expression is unity, and one obtains the lower bound on $V_{tb}$ quoted above. However, if there are more than three generations, there is no bound on $V_{tb}$ from this measurement; one only learns that $V_{tb} > V_{ts}, V_{td}$. This point is understood instantly by beauty physicists; the fact that the bottom quark decays predominantly to charm does not determine $V_{cb}$, only that $V_{cb} > V_{ub}$.

### 1.1 Indirect information on $V_{tb}$

Before we proceed to the direct measurement of $V_{tb}$, let’s explore what we know about it indirectly. Consider first precision electroweak measurements, such as the $\rho$ parameter,

$$\rho \equiv \frac{M_W^2}{M_Z^2 \cos^2 \theta_W} = 1 + \frac{3G_F}{8\sqrt{2}\pi} m_t^2 \frac{m_t}{m_W^2}.$$ \hfill (2)

where the last term comes from the loop diagram in Fig. 1(a). The measurement of this parameter told us that the top quark mass must be less than about 200 GeV long before it was discovered. If there were a fourth generation, then it would also contribute to the $\rho$ parameter, via the loop diagram in Fig. 1(b). Recall that the loop correction to the $\rho$ parameter depends on the mass splitting within an SU(2) doublet. Since the $tb$ doublet already contributes as much as is necessary to yield a $\rho$ parameter consistent with experiment, the mass splitting in the fourth generation must be small, $m_t \approx m_{t'}$. Allowing for mixing between the third and fourth generations, Eq. (2) becomes

$$\rho \approx 1 + \frac{3G_F}{8\sqrt{2}\pi^2} \left[ m_t^2 |V_{tb}|^2 + m_{t'}^2 |V_{t'b}|^2 \right].$$ \hfill (3)

\textsuperscript{b}We know from LEP that if there is a fourth generation, the associated neutrino is not light; however, this does not rule out a fourth generation.

\textsuperscript{c}Furthermore, if there are more than three generations, there is no lower bound on $V_{ts}$ or $V_{cd}$.\hfill (3)
The consequences of Eq. (3) are laid bare if we invoke four-generation unitarity to write
\[ |V_{tb}|^2 = 1 - \epsilon^2, \quad |V_{t'b}|^2 = \epsilon^2. \]
We then obtain
\[ \rho \approx 1 + \frac{3G_F}{8\sqrt{2}\pi^2} \left[ m_t^2 + \epsilon^2 (m_t'^2 - m_t^2) \right] \]  
\[ (4) \]
What we learn from Eq. (4) is that either $\epsilon^2$ is small, which corresponds to $V_{tb} \approx 1$, or $m_t' \approx m_t$. The lesson is that in order for $V_{tb}$ to deviate significantly from unity, there must be a fourth generation of quarks whose mass is not far above that of the top quark.

As beauty physicists know well, the top quark also contributes indirectly to flavor-changing neutral current processes such as $b \to s\gamma$, as shown in Fig. 2(a). The observed rate for $b \to s\gamma$ is consistent with $V_{tb} \approx 1$. If there were a fourth generation, it would also contribute to $b \to s\gamma$, via Fig. 2(b).

The total contribution to the amplitude from the third and fourth generation is proportional to
\[ A \sim F_2(\frac{m_t'^2}{M_W^2})V_{tb}V_{ts}^* + F_2(\frac{m_t^2}{M_W^2})V_{t'b}V_{t's}^* \]  
\[ (5) \]
where $F_2(x)$ is the Inami-Lim function. Using the four-generation unitarity relation $V_{tb}V_{ts}^* + V_{t'b}V_{t's}^* + V_{cb}V_{cs}^* + V_{ub}V_{us}^* = 0$, and neglecting the last term in this relation (known to be small), we can rewrite Eq. (3) as
\[ A \sim -F_2(\frac{m_t'^2}{M_W^2})V_{cb}V_{cs}^* + [F_2(\frac{m_t^2}{M_W^2}) - F_2(\frac{m_t'^2}{M_W^2})]V_{tb}V_{t's}^*. \]  
\[ (6) \]
The first term above is all that is needed to produce a rate for $b \to s\gamma$ consistent with experiment. The second term is negligible if either $m_t' \approx m_t$, or $V_{t'b}V_{t's}^*$

\[ ^a \text{I am assuming } V_{ts}, V_{td} \ll 1 \text{ throughout this argument, for simplicity. Relaxing this assumption does not change the conclusion of the argument.} \]
is small, regardless of the value of $V_{tb}$. Thus this process cannot constrain $V_{tb}$ if there are more than three generations.

We conclude that, although $V_{tb} \approx 1$ is consistent with all known phenomenology, the case is not airtight. We should therefore measure it directly, and keep our eyes and minds open.

1.2 Direct measurement of $V_{tb}$

The best way to measure $V_{tb}$ at hadron colliders is via single-top-quark production. There are two separate processes: quark-antiquark annihilation shown in Fig. 3(a), and $W$-gluon fusion shown in Fig. 3(b). Both proceed via the weak interaction, with a cross section proportional to $|V_{tb}|^2$.

A strategy to extract $V_{tb}$ from single-top-quark production is as follows. One measures the cross section times the branching ratio of $t \rightarrow bW$,

$$\sigma = \sigma(t\bar{b}) BR(t \rightarrow bW).$$  \hspace{1cm} (7)

To determine $BR(t \rightarrow bW)$, one uses the measured cross section for $t\bar{t}$ production and decay,

$$\sigma = \sigma(t\bar{t})[BR(t \rightarrow bW)]^2.$$

One extracts $BR(t \rightarrow bW)$ from this measurement, and inserts it into Eq. (7) to obtain $\sigma(t\bar{b})$, which is proportional to $|V_{tb}|^2$. This procedure requires an accurate theoretical calculation of both $\sigma(t\bar{b})$ and $\sigma(t\bar{t})$. The former has been calculated to next-to-leading-order in the strong interaction for both quark-antiquark annihilation and $W$-gluon fusion. The latter has also been calculated to next-to-leading-order in the strong interaction, and is known to $\pm 10\%$. The next-to-next-to-leading-order correction for all three processes would be desirable.

\textsuperscript{12} There are attempts to go beyond next-to-leading order by summing the effects of soft gluons.\textsuperscript{12,13}
The quark-antiquark process has the advantage that the quark distribution functions are well known, while the $W$-gluon fusion process suffers from the uncertainty in the gluon distribution function. Furthermore, one can exploit the similarity of the quark-antiquark-annihilation process to Drell-Yan to measure the incident quark-antiquark flux. However, $W$-gluon fusion has the advantage of a larger cross section at the Tevatron, and a much larger cross section at the LHC; it is doubtful that the quark-antiquark-annihilation process will be observable above backgrounds at the LHC. Thus the two processes are complementary as measures of $V_{tb}$.

Both quark-antiquark annihilation and $W$-gluon fusion should be observed in Run II at the Tevatron, and yield a measurement of $V_{tb}$ with an accuracy of $\pm 0.1$. Further running at the Tevatron (30 fb$^{-1}$) could reduce the uncertainty to $\pm 0.05$. This is comparable to the accuracy on $V_{tb}$ expected from a measurement of the top-quark width at the $t\bar{t}$ threshold in $e^+e^-$ and $\mu^+\mu^-$ colliders. The accuracy with which $V_{tb}$ can be measured at the LHC via $W$-gluon fusion is limited only by systematics, most notably the uncertainty in the gluon distribution function.

1.3 More on $W$-gluon fusion

Let’s look more closely at the calculation of the cross section for single-top-quark production via $W$-gluon fusion. The leading-order cross section, shown in Fig. 4(a), is proportional to $\alpha_s \ln(m_t^2/m_b^2)$. The large logarithm arises from the region of collinear $b\bar{b}$ production from the initial gluon. Another power of this large logarithm appears at every order in perturbation theory via the emission of a collinear gluon from the internal $b$ quark, as shown in Fig. 4(b). The resulting expansion parameter is thus $\alpha_s \ln(m_t^2/m_b^2)$, which yields a series
which is much less convergent than an expansion in $\alpha_s$ alone.

Fortunately, a technology exists to sum this large logarithm to all orders in perturbation theory. The idea is to use the DGLAP equations to sum the collinear logarithms. In the process, one generates a perturbatively-derived $b$ distribution function, as shown in Fig. 4(c). The diagram in Fig. 4(c) becomes the leading-order diagram; it is of order $\alpha_s \ln(m_t^2/m_b^2)$, because the $b$ distribution function is intrinsically of this order. The non-collinear part of the $W$-gluon fusion process (Fig. 4(a)) is only of order $\alpha_s$, so it corresponds to a correction to the leading-order process (Fig. 4(c)) of order $1/\ln(m_t^2/m_b^2)$.

Virtual- and real-gluon emission corrections to Fig. 4(c) correspond to corrections of order $\alpha_s^2$.

Both beauty and charm should have perturbatively-derivable distribution functions, since $m_b, m_c >> \Lambda_{QCD}$. This can be tested via heavy-flavor production at HERA. For $Q^2 >> m_c^2$, the leading-order process for charm production is shown in Fig. 5, with a charm distribution function. The observed cross section should agree with that calculated via Fig. 5 (and radiative corrections to it), using a perturbatively-derived charm distribution function. If it does not agree, then one must appeal to “intrinsic charm”, which corresponds to a nonperturbative component to the charm distribution function.

2 $m_t$ and the $t\bar{t}$ threshold

Although the top quark was only recently discovered, the top-quark mass is already the second best-known quark mass (as a percentage of its mass). The top-quark mass is obtained from the invariant-mass distribution of its decay
products, a $W$ and a $b$ jet. The peak in this distribution corresponds closely with the top-quark pole mass. The CDF/D0 average is $m_t = 175.6 \pm 5.5$ GeV.  

\begin{equation}
 m_t = 175.6 \pm 5.5 \text{ GeV}.
\end{equation}

This corresponds to an $\overline{\text{MS}}$ mass of 

\begin{equation}
 \overline{m}_t(m_t) = 166.5 \pm 5.5 \text{ GeV}.
\end{equation}

In contrast, the bottom and charm masses are obtained from the $\Upsilon$ and $\Psi$ spectra. Since these spectra involve nonperturbative QCD, the most reliable method to obtain the quark mass is via lattice QCD. This yields the $b$-quark mass $m_b$:

\begin{equation}
 \overline{m}_b(m_b) = 4.0 \pm 0.1 \text{ GeV},
\end{equation}

which is the best-known quark mass.

Amongst the quarks, as well as the leptons, the top quark is unique; it is the only fermion whose mass lies near the electroweak scale, which is the natural value for fermion masses in the standard model. Fermions acquire their mass via their Yukawa coupling to the Higgs field, $y_f$:

\begin{equation}
 m_f = y_f \frac{v}{\sqrt{2}}
\end{equation}

where $v \approx 250$ GeV is the vacuum-expectation value of the Higgs field. The Yukawa coupling of the top quark is near unity, a natural value, while all other fermions have a Yukawa coupling much less than unity. It therefore seems likely that, if we are ever to have a theory of fermion masses, the top-quark mass will be the first to be understood. This provides a motivation for measuring its mass as precisely as possible.
The electroweak gauge couplings have been measured to a precision of about 0.1%. This provides a goal for the accuracy of the top-quark mass measurement, roughly 200 MeV. Such an accuracy cannot be achieved at hadron colliders; the best prospect is to make a precision measurement of the $t\bar{t}$ threshold in $e^+e^-$ or $\mu^+\mu^-$ colliders. This is the analogue of the extraction of the bottom and charm masses from the $\Upsilon$ and $\Psi$ spectra. However, due to the short top-quark lifetime, topononium states do not have time to form, so the $t\bar{t}$ threshold is rather different from that of bottom and charm.

2.1 Toponium spectroscopy

It is amusing to consider what toponium spectroscopy would be like if the top quark lived long enough to form toponium states. The ground state would be very small, with a Bohr radius ($C_F = 4/3$ is a color factor)

$$a = (C_F \frac{m_t}{2} \alpha_s)^{-1} \approx (12 \text{ GeV})^{-1},$$

which is much less than the characteristic distance scale of nonperturbative QCD, $\Lambda_{\text{QCD}}^{-1} \approx (200 \text{ MeV})^{-1}$. The ground state, as well as the first few excited states, would be governed almost entirely by perturbative QCD, which gives rise to a potential

$$V(r) = -C_F \frac{\alpha_s(1/r)}{r},$$

with a coupling that depends on the distance between the quark and the antiquark. If we ignore the running of the coupling, we can use this Coulomb potential to estimate the number of toponium levels there would be below $T\bar{T}$ threshold ($T = t\bar{q}$ is a top meson). This threshold is attained when the size of the toponium state is of order $\Lambda_{\text{QCD}}^{-1}$. Recall that the size of a Coulomb bound state is approximately $n^2a$, where $n$ is the principal quantum number. Using Eq. (13) we obtain

$$n_{\text{max}} \approx 0.8 \left( \frac{m_t}{\Lambda_{\text{QCD}}} \right)^{1/2}.$$  

\(^{15}\)

The constant of proportionality in Eq. (15) was obtained by fitting this formula to the $\Upsilon$ spectrum, which has $n_{\text{max}} \approx 4$. We find that for toponium, the

\(^{9}\)An example of a theory of fermion masses in which a precision of 1 GeV on the top-quark mass is sufficient is SO(10) grand unification with Yukawa-coupling unification.

\(^{h}\)The result that $n_{\text{max}} \sim m_t^{1/2}$ is actually true for an arbitrary potential.

\(^{i}\)This includes the $\Upsilon(4s)$, which is only slightly above the $B\bar{B}$ threshold. Eq. (15) also works well for the $\Psi$ spectrum, which has $n_{\text{max}} \approx 2$. 

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number of principal quantum numbers below threshold would be

\[ n_{\text{max}} \approx 24. \]  \hspace{1cm} (16)

Taking into account angular excitations, spin-orbit and spin-spin interactions, this corresponds to \( 2n_{\text{max}}^2 \approx 1152 \) distinct energy levels below threshold.

2.2 \( t\bar{t} \) threshold and the top-quark width

In reality, toponium states do not have time to form before the top quark decays. The formation time can be estimated by the size of the toponium state divided by the velocity of the quark and antiquark, which in a Coulomb bound state is approximately \( C_F \alpha_s \):

\[ t_{\text{form}} \approx \frac{n^2 a}{C_F \alpha_s} \approx n^2 (1.6 \text{ GeV})^{-1}. \]  \hspace{1cm} (17)

[The formation time can also be estimated via \( t_{\text{form}} \approx 1/|E_n| \), where\]

\[ E_n = -\frac{1}{2} \frac{m_t C_F \alpha_s^2}{n^2} \]  \hspace{1cm} (18)

is the energy of the \( n^{\text{th}} \) level (relative to \( 2m_t \)); this also yields Eq. (14).] The formation time is greater than the toponium lifetime, \( (2\Gamma)^{-1} \approx (3 \text{ GeV})^{-1} \) (half the top-quark lifetime, since either the \( t \) or \( \bar{t} \) can decay), for all but perhaps the ground state. As a consequence, the \( t\bar{t} \) threshold is a smooth continuum, with no sharp resonances, as shown in Fig. 6, curve (c). The 1s resonance is pronounced if we make the top quark width half its standard-model value, as shown in curve (a) of Fig. 6.

The top-quark width plays a more important role than just smoothing out the \( t\bar{t} \) threshold, however; it acts as an infrared cutoff, eliminating the contribution to the threshold from nonperturbative QCD. \[ \] One can see this in a heuristic way as follows. Imagine that the top quark is stable, and consider the wave function of the \( n^{\text{th}} \) state, which falls off exponentially as

\[ \psi_n \sim e^{-r/n a} \sim e^{-r \sqrt{-m_t E_n}} \]  \hspace{1cm} (19)

where I have used Eqs. (13) and (18) to express the argument of the exponential in terms of the energy. In reality there is not sufficient time for the bound states to form; one can implement this information by saying that the bound states form in the complex energy plane at \( E_n - i(2\Gamma)/2 \). To find the impact of these resonances on physics at real values of \( E \), we let \( E_n \rightarrow E_n + i\Gamma \) in Eq. (19), to obtain

\[ \psi_n \sim e^{-r \sqrt{-m_t (E_n + i\Gamma)}} \]  \hspace{1cm} (20)
In the absence of the width, the exponential suppression of the wave function
at long distances is reduced as one approaches threshold, $E_n \to 0$, and the
toponium state has a large nonperturbative contribution from distances of
order $\Lambda_{\text{QCD}}^{-1}$ and greater. However, with the width present, an exponential
suppression remains even as one approaches threshold:

$$|\psi_n| \xrightarrow{E_n \to 0} e^{-r \sqrt{m_t \Gamma}/2}. \quad (21)$$

The width suppresses the wave function at distances greater than $1/\sqrt{m_t \Gamma} \approx
(16 \text{ GeV})^{-1}$, which is much less than $\Lambda_{\text{QCD}}^{-1} \approx (200 \text{ MeV})^{-1}$. Thus the $t\bar{t}$
threshold has a negligible contribution from nonperturbative QCD. The thresh-
old line shape, as shown in Fig. 6, can be calculated entirely within perturba-
tion theory, limited in accuracy only by our strength to carry out higher-order
calculations.

2.3 Top-quark pole mass

Because the $b$ quark is confined into hadrons, we are used to the notion that the
$b$-quark pole mass is unphysical. For example, one could attempt to identify
the $b$-quark pole mass as the mass of a $B$ meson minus the binding energy of
the light quark, but this binding energy is an ambiguous concept, due to con-
finement. Since the binding energy is of order $\Lambda_{\text{QCD}}$, one expects an ambiguity
in the $b$-quark pole mass proportional to $\Lambda_{\text{QCD}}$.

The above argument can be made more rigorous by studying the one-loop
contribution to the $b$-quark propagator, as shown in Fig. 7. The region of soft
gluon momentum inevitably involves nonperturbative QCD. It has been shown
that this prevents an unambiguous definition of the $b$-quark pole mass, with an
ambiguity proportional to $\Lambda_{\text{QCD}}$.

Given that the width acts as an infrared cutoff in top physics, scree-
ning the effects on nonperturbative QCD, one might expect the width to allow an
unambiguous definition of the top-quark pole mass. This is not the case. The
width moves the perturbative top-quark pole into the complex energy plane,
but the position of the real part of the pole remains ambiguous by an amount
proportional to $\Lambda_{\text{QCD}}$. The width does not remove the contribution of soft
gluons to the top-quark pole mass.

For example, consider the measurement of the top-quark mass at hadron
colliders by reconstructing the invariant mass of the decay products, a $W$ and
a $b$ quark. In reality, the $b$ quark manifests itself as a jet of colorless hadrons.
Since the top quark is colored, at least one of the quarks in the hadrons that
make up the $b$ jet must not be a decay product of the top quark. Since there is
no unambiguous way of subtracting the binding energy of this quark from the
$b$ jet, the $Wb$ invariant mass is ambiguous by an amount proportional to $\Lambda_{\text{QCD}}$. Thus
the top-quark pole mass is ambiguous by an amount proportional to $\Lambda_{\text{QCD}}$
despite the fact that the top quark does not form hadrons.

The moral is that if one hopes to measure the top-quark mass to an accu-
racracy of 200 MeV, one must use a short-distance mass, such as the $\overline{\text{MS}}$ mass,
in the calculation. There is no fundamental roadblock to measuring this mass
to arbitrary precision from the $t\bar{t}$ threshold.

3 Conclusions

Because top-quark physics is free from the effects of nonperturbative QCD, we
already have precision top-quark theory. What is needed is precision top-quark
experiments. We await a measurement of $V_{tb}$ to $\pm 0.1$ from single-top-quark production in Run II at the Tevatron. Additional running ($30 \text{ fb}^{-1}$) could reduce the uncertainty to $\pm 0.05$. The accuracy on $V_{tb}$ from $W$-gluon fusion at the LHC will be limited only by systematic errors, in particular the uncertainty in the gluon distribution function. The top-quark MS mass can be measured to 200 MeV or even better at $e^+e^-$ and $\mu^+\mu^-$ colliders running at the $t\bar{t}$ threshold. By making precision studies of the top quark we hope not only to measure the standard-model parameters accurately, but to uncover physics beyond the standard model.

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