TOP PHYSICS AT LHC

CONVENERS: E. Reya and P. M. Zerwas


* LHC Workshop, Aachen, Oct. 1990
The search for new quarks and the exploration of their properties has been a most important task at hadron colliders in the past. Proton-(anti)proton colliders may even be the only machines in which top quarks can be produced in this decade, thus completing the fermion spectrum of the Standard Model. The present lower limit [1] of 89 GeV for the top quark mass leaves open only a small window for LEP200. Estimates based on precision measurements of the electroweak parameters indicate a value between 100 and 220 GeV for the top quark mass [2]. Part of this range or eventually even the entire range can be covered at the Tevatron. However, if the mass were close to 200 GeV or beyond, it may need the LHC to discover this long awaited particle. And only a high-energy collider can provide the large number of top quark events [order $10^7$] necessary to measure the top quark mass accurately within a few GeV and to determine the decay properties.

The main production mechanisms for top quarks in proton-proton collisions are gluon-gluon and quark-antiquark fusion (Fig.1a) [3]

$$gg \rightarrow t\bar{t}$$

for quark masses up to $\sim 250$ GeV. Only if top quarks were still heavier, $W$-gluon fusion would become the dominant production process (Fig.1a) [4]

$$Wg \rightarrow t\bar{b}$$

Estimates of these cross sections are shown in Fig.1b for various parametrizations and models of the parton densities [5,6]. Depending on the mass, a number of $10^5$ to $10^6$ top quarks can be created at a modest integrated luminosity of $\mathcal{L} = 10^3$ pb$^{-1}$.

---

*E. Reya and P. M. Zerwas*
As demonstrated in Fig.1c, the discovery limit [defined by 1,000 events] for any sort of heavy quarks in gg/qg fusion amounts to $m_Q \approx 1$ TeV for an integrated luminosity of $10^4$ pb$^{-1}$.

Heavy top quarks decay semiweakly into bottom quarks plus W bosons [7]

$$t \rightarrow b + W^+$$

Their width, however, remains rather narrow $\lesssim 3$ GeV up to top quark masses $\lesssim 200$ GeV. The branching ratios of the semileptonic decay modes are thus the same as for the leptonic decay modes of the W boson,

$$BR(t \rightarrow \ell^+) = 1/9$$

and a fraction of about 20% decay promptly into isolated muons and electrons. Additional leptons that are emitted in semileptonic $b$-quark decays [including $b \rightarrow c \rightarrow l$ cascades], are non-isolated and associated with hadron jets.

Top-antitop quark pairs decay into 0-lepton, 1-lepton, 2-lepton, ... final states [8], the leptons being either isolated or non-isolated decay products of $b$-quarks. The large QCD jet backgrounds render purely hadronic top quark decays unetectable. Single-lepton top events are difficult to isolate because of the large $W + n$-jet background [9]. Applying a set of stringent cuts on the transverse momenta of leptons and jets, the signal can nevertheless be extracted at an S/B level of order 2. Demanding further the invariant masses of dijet pairs to be in the $W$ mass range or tagging the $b$ jets reduces the background very efficiently. A very clean sample can easily be obtained from events with at least 2 [isolated] leptons. Since the total sample of top quarks is very large at the LHC, the requirement of two leptonic W decays does not impede a high-statistics evaluation of the signal.

The top quark mass can be determined by several methods: (i) The strong dependence of the fusion cross section on the quark mass (Fig.1b) can in principle be exploited to get a first estimate of the top mass. Uncertainties are introduced through higher-order QCD corrections, parton densities and fragmentation functions. After applying selection cuts on the lepton final states, the mass dependence is less pronounced, rendering this method less suitable for an accurate determination of the top mass at high energies; a limit of $\pm 15$ GeV may finally be achieved. (ii) A model independent measurement of the top quark mass within $\pm 8$ GeV can be carried out by studying the invariant mass distribution of the jets recoiling against the lepton in 1l events. (iii) Single-lepton events in a top sample, that is sufficiently large to allow for tight cuts, will enable us to measure the top quark mass directly by a second method. Reconstructing the longitudinal neutrino momentum from the $W$ mass constraint on the $(b\ell\nu)$ pair, the top mass can be read off the peak in the $W + b$-jet mass distribution [10]. (iv) The best experimental value of the top quark mass can be derived from the dilepton mass distribution in 3-lepton events when an isolated lepton is paired with an oppositely charged lepton belonging to a $b$-quark jet [11]. The main systematic error source is the uncertainty in the $b$-quark fragmentation. As a consequence of environmental independence, however, the fragmentation parameters are nearly the same as those in $Z^0$ decays because the $b$-quark energies are quite similar in both cases. Based on the fragmentation parameters measured at LEP, the error on the top mass may be reduced to $\Delta m_t \approx \pm 5$ GeV.

The second observable, that can also be extracted in a model independent way, is the semileptonic branching ratio. Comparing the yield of 1-lepton events with 2-lepton events, this branching ratio can be evaluated with an accuracy of $\Delta BR_t \approx \pm 5\%$, without any reference to theoretical predictions of the production cross section. No method is available for measuring the top quark width in such a machine. Even though the possibility may be rather remote that the weak ($tb$) charged current be different from the usual V-A form of the Standard Model, the chirality of this current can nevertheless be investigated by measuring the final state $(tb)$ invariant mass distribution [12]. Fixing the chirality of the $(t\nu)$ lepton current to the standard V-A form, this distribution is sensitive to the chirality of the weak ($tb$) charged quark current.

The radiative QCD corrections to the total cross section and to the inclusive single-top distribution of the fusion processes have been completed to second order [13,14]. Measurements of the production cross sections therefore allow stringent QCD tests. In particular, the extrapolation of quark and gluon densities from the present energy range to the high LHC energies can be scrutinized thoroughly.

Top physics is completely predicted in the Standard Model once the top quark mass, the presently only unknown parameter of this particle, is fixed. This area is therefore an ideal test ground for possible phenomena beyond the Standard Model.

(i) It is very difficult to detect charged Higgs particles in hadron colliders that can be created directly only by the electroweak Drell-Yan mechanism. However, a good fraction of top quarks could decay to charged Higgs bosons [7,15]

$$t \rightarrow b + H^+$$

if $m_H < m_t - m_b$. In the minimal supersymmetric extension of the Standard Model, the charged Higgs particles decay predominantly into $\tau$ leptons [apart from possible light scalar decays $H^+ \rightarrow W^+ h^0$]. As a result, deviations from the universal semileptonic branching ratios would be observed in top decays. Depending on the $t$ mass, charged Higgs particles in the mass range between 100 and 150 GeV could be traced this way. This range can in principle be extended by associated production of top and Higgs particles in gb fusion [16].

(ii) If the supersymmetric partner $\tilde{t}$ of $t$ were lighter than the heavy top quark [17,18], this would manifest itself in lepton spectra softer than in standard decays and it would lead to an increase of the missing transverse energy. For example, if top quarks decay through the chain

$$t \rightarrow \tilde{t} + \tilde{\nu}$$

$$\rightarrow b + W^\pm$$

$$\rightarrow W + \tilde{\nu}$$

$$\rightarrow \mu + \nu_\tau$$

a large fraction of energy is deposited in the invisible neutralinos ($\approx$ photinos) $\tilde{\nu}$ and the neutrino.

(iii) Mixing [19] between heavy quark states of order $m_t/m_b/s^2$ that would be suppressed for light quarks, could manifest itself in NC decays of top quarks

$$t \rightarrow c + Z$$
complementing the standard CC decays to b quarks. Since leptonic Z decays provide a striking signature, this mode would be easy to detect.

(iv) Other extensions of the Standard Model that involve a fourth family [20], necessarily incorporate heavy neutrinos. In this case two kinematical possibilities could be realized: For $m_N < m_t + m_W$ the signatures of $b 	o W W b$ decays would be very similar to the standard top decay $t \to W b$. If $m_N > m_t + m_W$, the decays of pair produced $b$ quarks in $gg/qq$ fusion

$$b\bar{b} \to WWb \to WWWbb$$

give rise to large samples of four-W events.

(v) Extensions of the families themselves are predicted in GUTs like E(6) where the chiral quark spectrum is supplemented by isoscalar down-type quarks [21]. These exceptional particles may decay either through SUSY chains or through mixing with ordinary light quarks, inducing spectacular FCNC decay modes in addition to the usual CC decays [22]

$$D \to d + Z$$

$$\to u + W^-$$

Based on 1- and 2-lepton events accompanied by several hadronic jets, discovery limits of

$$m(D) \leq 400 \text{ GeV}$$

are predicted for the LHC collider at an integrated luminosity of $L = 10^4 pb^{-1}$.

The report of the Top Working Group is split into three parts:

1.) Theoretical basis [following this summary]
   (E. Reya, P. Zerwas, W. Hollik, V. Khoze, R.J.N. Phillips, F. Berends, J. Zunfi)

2.) Top search: signal vs. background
   (F. Cavanna, D. Denegri, T. Rodrigo)

3.) Top mass and decays
   (L. Fayard, M. Felcini, H. Reithler, G. Unal)

Acknowledgements: We thank many colleagues for discussions and for their contributions to various topics of this investigation, in particular to K. Eggert (CERN), D. Froidevaux (CERN), J. Ohnenius (Florida State Univ.), J. Vidal (Univ. Valencia), W. Vogelsang (Univ. Dortmund) and D. Zeppenfeld (Univ. Wisconsin).

REFERENCES

[1] CDF Collaboration, 25th Rencontre de Moriond, Les Arcs 1990 (pres. by K. Sliwa); for earlier analyses see:
   C. Albajar et al. (UA1 Coll.), Z. Phys. C48 (1990) 1;
   T. Akesson et al. (UA2 Coll.), Z. Phys. C46 (1990) 179;

   D. Haidt, DESY 89-073;
   P. Langacker, UPR-0435 T (1990);


1. Production Mechanisms for Top Quarks

The main production mechanisms for top quarks at high energy proton-proton colliders are the $gg \rightarrow t\bar{t}$ and $q\bar{q} \rightarrow t\bar{t}$ fusion processes. As demonstrated in the previous section, only for top quark masses beyond $\sim 250$ GeV semileptonic $WW$ fusion to $t\bar{t}$ becomes dominant. In the following subsections these production mechanisms and the corresponding characteristics of the final states are summarized, with the strong and electroweak radiative corrections for the $gg$ and $q\bar{q}$ fusion processes included.

1.1 Hadronic Fusion Mechanisms

Before entering a more detailed discussion of higher order QCD contributions let us first discuss the dominant Born terms for the total cross sections and for the characteristics of the final states. The partonic leading order (LO) cross sections for $gg \rightarrow t\bar{t}$ and $q\bar{q} \rightarrow t\bar{t}$ are known to be [1]

$$\sigma_{gg}(s) = \frac{\pi^2}{3s} \left[ \ln \left( 1 + \beta \rho \right) - \beta \left( \frac{31}{4} + \frac{3}{16} \rho \right) \right]$$

$$\sigma_{q\bar{q}}(s) = \frac{8\pi^2}{27s} \beta \left[ 1 + \frac{\rho}{1} \right]$$

(1.1)

(1.2)

with $\rho = 4m_t^2/s$ and $\beta = \sqrt{1 - \rho}$ being the velocity of the $t$ quarks in the $t\bar{t}$ cm frame with invariant energy $s$. The arbitrary renormalization scale of the coupling will be chosen, for the time being, to be $\Lambda$; other choices will be discussed below. The total pp cross sections then follow by averaging these partonic cross sections over the $q\bar{q}$ and $gg$ luminosities in pp collisions. The two curves limiting the band for $\sigma(pp \rightarrow t\bar{t})$ in Fig.1b of the Introduction are based on the DFLM [2] and (LO) GRV [3] distributions. Gluon fusion is by far the dominant production mechanism, as evident from Table 1 and intuitively expected. With a total cross section between $10^2$ pb at $m_t \sim 300$ GeV and $10^4$ pb at $m_t \sim 100$ GeV, a few times $10^2$ up to $10^4$ top quarks will be produced even at a modest integrated luminosity of $\int L = 10^{56}$ pb$^{-1}$. This leaves ample room for necessary acceptance corrections in order to delineate a clean top signal. Furthermore, the average invariant energy $<\sqrt{s}>$ of the $t\bar{t}$ pair is just a few $m_t$ as can be read off the representative mass values in Table 1. The top production is central with a spread of $\Delta y = 2$, the rapidity values where the cross section has fallen to $1/2$ of its maximum. The distribution of the transverse momenta has its maximum value at $p_T^{\text{max}}(t) \sim m_t/2$.

More generally, the leading order LO and next-to-leading order NL inclusive hadronic production cross section of a heavy top-quark pair is generically given by [1,4-6]

$$d\sigma_{pp} = \sum_{i,j} d\sigma_{i,j} d\Phi_{i,j} F_i(x_i, \mu_i^2) F_j(x_j, \mu_j^2) d\phi_{i,j}(x_i, x_j, \xi, \ldots, m_{i,j}^2, \mu^2)$$

(1.3)

where $F_i$ are the densities of parton $i$ in the proton and $\mu_i$ is the factorization/renormalization scale. As already indicated above, the partonic cross sections $d\sigma_{i,j}$ consist in the dominant LO of the $O(\alpha_s^2)$ Born contributions $gg \rightarrow t\bar{t}$, $q\bar{q} \rightarrow t\bar{t}$ in (1.1) and (1.2); in a full higher-order (HO) calculation the next-to-leading $O(\alpha_s^3)$ $2 \rightarrow 3$ contributions $gg \rightarrow t\bar{t}q$, $q\bar{q} \rightarrow t\bar{t}q$ enter as well and the total HO contribution is the sum of both. Before turning to these cross sections let us first briefly discuss our choice for the parton distributions needed in eq.(1.3).

For a consistent theoretical calculation it is obvious from the above discussion that we need two distinct sets of parton distributions. For a LO calculation one needs distributions calculated using LO evolution equations together with the LO (1-loop) expression for $\alpha_s(\mu^2)$; for a HO calculation the distributions have to be calculated using the HO evolution equations with the HO (2-loop) expression for $\alpha_s(\mu^2)$ employing the same subtraction scheme as used for calculating the cross sections of the subprocesses. Needless to say that, for a perturbative calculation to make sense and be reliable, the HO results should not differ too much from the LO ones, even taking into account the various ambiguities inherent in eq.(1.3) such as the choice of $\Lambda$, $m_t$, $\Lambda$, $F(x, \mu^2)$ and a specific renormalization scheme. As a typical example of 'conventional' parton distributions we use the LO and HO parametrization of DFLM [2] where $\Lambda_{DFLM} = 153$ MeV for $\alpha_s(\mu^2)$ and in HO $\Lambda_{DFLM} = 100$ MeV for $\alpha_s(\mu^2)$, respectively. [Note that $\Lambda_{DFLM} = 153$ MeV corresponds to $\Lambda_{DFLM} = 200$ MeV, and $\Lambda_{DFLM} = 100$ MeV is related to $\Lambda_{DFLM} = 160$ MeV which follows from the continuity constraint $\alpha_s^{(1)}(m_t) = \alpha_s^{(1)}(m_t)$ across the mass boundary $m_b = m_t$.] Here, the input valence gluon and sea distributions at $Q^2 = 10$ GeV$^2$ are extracted from deep inelastic data (where $x > 0.01$) using the characteristic 'conventional' Regge ansatz $xG(x, Q^2) \sim \text{const.}$ and $xq(x, Q^2) \sim \text{const.}$ as $x \rightarrow 0$. As an extreme alternative we use the 'radiative' GRV distributions [3] where only a valence-like input $G(x, \mu_0^2) = \alpha_s^{\text{eff}}(x, \mu_0^2)$, with $V = x_b + d_b$, and $q(x, \mu_0^2) \equiv 0$ is used at $\mu = \Lambda_Q$. The gluon and sea densities at $Q > \mu$ are uniquely generated radiatively by the LO and HO evolution equations with only a single parameter $\eta$ which can be extracted from experiment ($\eta \simeq 2/3$). The resulting distributions are then consistent with all present data, even down to very small values of $Q^2$, but differ from conventional parametrizations for $x \lessgtr 0.01$ by being steeper as $x \rightarrow 0$, i.e. $xG(x, Q^2) \sim x^{-\delta}$ with $\delta > 0$, and similarly for $q(x)$.

The expressions for the fully inclusive HO partonic cross sections in (1.3) are too lengthy for being recorded here but they can be found in Refs.[4,5]. The integrated expressions for the total cross sections, however, have a much simpler form:

$$\sigma_{i,j}(x, m_{i,j}^2, \mu_j^2) = \frac{\alpha_s^2(\mu_j^2)}{m_i^2} \left[ \sigma_{i,j}(0) + 4\alpha_s(\mu_j^2) \left( f_{i,j}^{(0)}(0) + f_{i,j}^{(1)}(0) \ln \frac{\mu_j^2}{m_i^2} \right) \right]$$

(1.4)

where $i = x_1 \ldots x_s$ and the dominant LO contributions $f_{i,j}^{(0)}(0)$ are given in eqs.(1.1) and (1.2),

Table 1

| $m_{top}$ [GeV] | $\sigma_{top}/\sigma_{tot}$ | $\sqrt{s}(t\bar{t})$ | $|G|^2$ |
|-----------------|-----------------------------|---------------------|------|
| 100             | 0.944                       | 311                 |
| 125             | 0.931                       | 383                 |
| 150             | 0.924                       | 454                 |
| 200             | 0.903                       | 593                 |
| 250             | 0.882                       | 729                 |
| 300             | 0.861                       | 862                 |

*E. Reya and P. M. Zerwas

**
and in addition \( f^{(0)}_{t} = f^{(0)}_{\bar{t}} = 0 \). The subleading HO expressions of \( f^{(1)}_{t} \) and \( f^{(1)}_{\bar{t}} \) are given in Refs.\cite{4,5}. In both HO calculations \cite{4,5} the heavy quark contributions have been calculated within the on-shell renormalization scheme with a fixed "physical" mass \( m_q \). [This is the most natural definition for the subsequent top decay process]. These HO results have then to be used in conjunction with the running coupling \( \alpha_s(\mu^2) \) and together with light parton densities evolved by employing the 2-loop MS evolution equations. The mass factorization scale \( \mu \) should be chosen such as to minimize the HO contributions with respect to the LO prediction. To estimate the scale dependence of our predictions we adopt the extreme choices of \( \mu^2 = m^2 \) and \( \mu^2 = m^2 + p_T^2 \) where the second choice has to be inserted in course in eq.(1.3) before integration over \( p_T \).

The predictions for the total cross section as a function of \( m_t \) are shown in Fig.1 for various c.m. energies \( \sqrt{s} \). In going up from the present Tevatron energy (1.8 TeV) to LHC (16 TeV) one gains more than two orders of magnitude in rate, depending on the value of \( m_t \), whereas the subsequent increase to SSC energies (40 TeV) is much less spectacular and for top searches of minor importance. The subdominant \( 2 \to 3 \) contributions add up to less than 10\% with respect to the dominant LO results which gives us strong confidence in the these purely perturbative QCD predictions. The theoretical uncertainties of the predictions at LHC due to different LO and/or HO parton distributions are about \( \pm 10\% \) as shown in Fig.2 and the scale ambiguity \( \mu = m_t \to m_T, \sqrt{m_T^2 + p_T^2} \) amounts to an additional \( \pm 15\% \) variation. Note that the "\( K \) factor", defined formally by the higher order corrections to the LO parton cross section but the parton distributions and \( \alpha_s \) kept fixed, amounts to an \( \approx 50\% \) correction of the Born terms. For completeness the LO and HO predictions are displayed in Fig.3 for the rapidity distribution \( d\sigma/dy \). Most of the top-events are produced centrally, concentrated in the interval \( |y| \leq 2.5 \).

It is also of interest to study separate, physically distinct components of the full HO \( \alpha^3_s \) results \cite{7}. Whereas the initial (final) state bremsstrahlung (ISGB) processes, illustrated for the gluon initiated reactions in Fig.4, dominate around threshold \( (\sqrt{s} \geq 2 m_t \text{ or } p_T < m_T) \), the gluon splitting (GS) and flavor excitation (FE) contributions become equally important for \( \sqrt{s} \gg 2 m_t \) or \( p_T > m_T \). \cite{5,7}. These various contributions to \( d\sigma/dp_T \) are shown in Fig.5. It is evident again that the LO predictions describe remarkably well these HO results in shape as well as in magnitude. The difference between different parton distributions is also marginal.

---

**Fig.1.** Production cross section for top pairs in \( pp \) and \( pp \) colliders: Tevatron (1.8 TeV); Tevatron II (3.6 TeV); LHC (16 TeV); SSC (40 TeV).

**Fig.2.** Predictions of top quark cross sections for two different parametrizations of parton densities \cite{2,3} using \( \mu^2 = m_T^2 + p_T^2 \). The vertical bars indicate the ambiguity of the HO predictions due to using \( \mu = m_T \) or \( 2m_T \), and using also the largest possible value of \( \Lambda \) for the DFLM distributions, \( \Lambda_{pt}^{MS} = 250 \text{ MeV} \) (which corresponds to \( \Lambda_{pt}^{MS} = 360 \text{ MeV} \)).
Near the $t\bar{t}$ threshold the cross sections are affected by resonance production and Coulomb rescattering forces [8,9,10]. These corrections can be estimated in a simplified potential picture. The driving 1-gluon exchange potential is attractive if the $t\bar{t}$ pair is in a color singlet state and repulsive in a color-octet state

$$
\sigma^{(1)}(gg\rightarrow t\bar{t}) = \frac{2}{5} \sigma_B(gg\rightarrow t\bar{t}) |\Psi^{(1)}|^2 \\
\sigma^{(8)}(gg\rightarrow t\bar{t}) = \frac{5}{7} \sigma_B(gg\rightarrow t\bar{t}) |\Psi^{(8)}|^2 \\
\sigma^{(8)}(q\bar{q}\rightarrow t\bar{t}) = \sigma_B(q\bar{q}\rightarrow t\bar{t}) |\Psi^{(8)}|^2
$$

(1.5)

The Coulombic attraction leads to a sharp rise of the cross section at the threshold in the singlet channel, even if no resonances can be formed anymore, since the phase space suppression of the Born term $\sigma_B \propto \beta_t$ is neutralized by the Coulomb enhancement of the wave function $|\Psi_t|^2 \propto \alpha_s/\beta_t$. In the octet channel, by contrast, the cross sections are strongly reduced by the Coulombic repulsion which effectively leads to an exponential fall-off of the cross sections $\sigma \propto \exp(-\pi \alpha_s/6\beta_t)$ at the threshold. Due to the averaging over parton luminosities [10] the effects are less spectacular in $pp$ than in $e^+e^-$ collisions.

In going beyond pure QCD corrections, let us first note that the lowest order $\alpha_s^2$ electroweak contributions, the Drell-Yan annihilation process (via $\gamma$ and $Z^0$ exchange), would only be of significance if $m_t < M_Z/2$, what has experimentally been ruled out by LEP. Here we examine the next order electroweak contributions $\mathcal{O}(\alpha_s^3)$ to both the $q\bar{q}\rightarrow t\bar{t}$ and $gg\rightarrow t\bar{t}$ subprocesses and we give quantitative results for the corrections to the parton cross sections [11]. These corrections become important in the LHC range for two reasons: the large $q^2 = s$ in the parton frame and the large Yukawa coupling of the (virtual) Higgs bosons to the top quark.

The loop diagrams to the parton processes under consideration are shown schematically in Fig.6. The complete set is gauge invariant, IR finite, and also UV finite without renormalization of the strong coupling constant (only mass renormalization in the $t$-channel top propagator is required). This is different for pure electroweak processes where the corresponding vertex corrections are not gauge invariant and also not UV finite without coupling constant renormalization. A large fraction of the corrections is due to the Higgs boson $H_t$ together with its unphysical components $\Phi^\pm, \chi$ in the $R_t$-gauge, that have Yukawa couplings $m_t/M_W$ to the top quark.

Fig.7 shows the relative correction to the $q\bar{q}\rightarrow t\bar{t}$ cross section at the parton level. Due to the initial state correction from $Z$ exchange they are flavor dependent, and the results of Fig.7 correspond to $q = u$. The difference to $d$-type quarks, however, is only marginal and not of numerical importance. The loop diagrams can be summarized in terms of an effective quark-gluon vertex with finite form factors $F_{V,A,M}(q^2)$,

$$
\gamma(p) = \gamma(p) \gamma_{\lambda V} + \gamma_{\mu} F_A + (p - \bar{p})_{\mu} F_M + \ldots
$$

(1.6)

*W. Hollik*
As a result of the interference with the Born diagram, \( F_A \) does effectively not contribute to the 1-loop cross section. The magnetic term is much smaller in size than \( F_V \) and, not too close to the threshold where the corrections are large, practically negligible. Hence, the correction is essentially a \( f \)-independent correction factor to the differential as well as the integrated cross section. The increasing influence of the Higgs boson with increasing top mass \( m_t \) is reflected by the broadening of the area covered by the variation of \( M_B \) between 42 and 1,000 GeV.

A similar situation is encountered for the \( gg \rightarrow t\bar{t} \) process. The box diagrams lead to a more complicated structure of the 1-loop terms and make the corrections depend also on the production angle. Fig. 8 displays the relative correction to the differential parton cross section for transverse top production (in the parton cms). The variation with \( p_t \), however, is rather smooth. The corrections become smaller in the forward direction and with higher top masses more sensitive to changes in \( p_t \).

Summarizing, except for a small region close to the production threshold, the corrections are always negative and can become sizeably large, in particular if the top is very heavy. After convoluting the cross sections of the subprocesses with the parton distributions, a reduction of the Born cross section at a level of 10 to 20% is expected for \( m_t \leq 200 \) GeV.

1.2 \( Wg \) fusion to \( b\bar{b} \)

The interaction radius in the QCD \( gg \) fusion process shrinks with rising energy so that the cross section \( \sigma(gg \rightarrow t\bar{t}) \sim \alpha_s^2/\bar{s} \) (mod. \( \log s \)) vanishes asymptotically. By contrast, the interaction radius in the weak fusion process \([12]\) is set by the Compton wave length of the W boson and therefore asymptotically non-zero, \( \sigma \sim G_F^2 m_W^2/2\pi \). Folding this subprocess with the quark-gluon luminosities, the fall-off of the total cross section \( \sigma(gg \rightarrow b\bar{b}) \) is less steep than for the QCD fusion processes. As a result, the \( Wg \) fusion process will become dominant at large top quark masses \( \geq 250 \) GeV, while being suppressed at masses of order 100 GeV by about two orders of magnitude. Nevertheless, as shown in detail in Fig. 1b of the Introduction, more than \( 10^6 \) top quarks will be produced by this mechanism for any mass value at an integrated luminosity of \( \mathcal{L} = 10^4 \) pb\(^{-1}\).

A close inspection of the diagrams in Fig. 1a reveals immediately that the by far dominant part of the cross section is due to \( b \) exchange, with the \( b \) quark being near its mass shell. Since the \( b \) quark is almost collinear to the incoming gluon, this cross section is logarithmically enhanced \( \sim \ln(m_t^2/m_b^2) \) over other mechanisms. This naturally suggests to approximate \([12]\) the process by the subprocess \( u + b \rightarrow d + t \) with the \( b \)-quark distribution generated perturbatively by gluon splitting employing massless evolution equations. The weak cross sections can be presented in a compact form,

\[
\sigma(ub \rightarrow dt) = \frac{G_F^2 m_W^2}{2\pi} \frac{\bar{s}m_t^2}{\bar{s} + m_b^2 - m_t^2},
\]

\[
\sigma(db \rightarrow ut) = \frac{G_F^2 m_W^2}{2\pi} \left[ \frac{1}{\bar{s}} \left( \frac{2m_b^2 - m_t^2}{\bar{s}} \right) - \frac{1}{\bar{s}} \left( \frac{2m_b^2 - m_t^2}{\bar{s}} \right) \log \frac{\bar{s} + m_b^2 - m_t^2}{m_b^2} \right],
\]

\[
\rightarrow \frac{G_F^2 m_W^2}{2\pi}.
\]

*D. Rein, P. M. Zerwas and J. Zunft*
and identically the same expressions for the C-conjugate reactions.

The characteristics of the final states are displayed in Table 2 for the subprocess

$$u + g \rightarrow d + t + \bar{b}$$

at a top quark mass of 150 GeV. The final state $d$ quark is emitted at a transverse momentum of $\sim m_W/2 \approx 40$ GeV that is characteristic for $W$ exchange processes. The rapidity distribution has its maximum value at $|y| \approx 3.5$. Both observables can be exploited to tag this type of reaction [13]. The spectator $\bar{b}$ quark, on the other hand, is characterized by a transverse momentum of $m_b \sim 5$ GeV and an average rapidity of $\sim \pm 1.6$ for events with positive/negative $y$. The top quark, finally, balances the transverse momentum of the $d$ quark. It is emitted in the forward direction [which is defined by the tagged $d$ quark]. Detailed distributions of the transverse momenta and the rapidities are shown in the experimental section. To illustrate the information that can be extracted from $d$ tagging, we present, in addition to the symmetrized distributions, "one-sided" $y$ values where the $W$ is assumed to belong to the proton travelling in $+z$ direction. [This case can approximately be realized by tagging the outgoing $d$ quark.]

Top quarks are created in $u + g$ collisions, anti-top quarks in $d + g$ collisions where the absorption of a $W^-$ transforms a $b$ quark to a $\bar{t}$ quark. The naive expectation from valence quark counting for the ratio of $t$ and $\bar{t}$ cross sections, $\sigma(u \rightarrow t) : \sigma(d \rightarrow \bar{t}) \sim 2 : 1$ is corroborated by a detailed analysis; in fact, the ratio turns out to be 2.1 for top quark masses of 150 GeV.

<table>
<thead>
<tr>
<th>$m_t$ (150 GeV)</th>
<th>Transv. momentum [GeV]</th>
<th>Rapidity</th>
<th>Rapidity (one-sided)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>do/dy = max</td>
<td></td>
<td>do/dy = max</td>
</tr>
<tr>
<td></td>
<td>etc.</td>
<td></td>
<td>etc.</td>
</tr>
<tr>
<td>$u + g \rightarrow d \bar{(t)}$</td>
<td>38</td>
<td>36</td>
<td>3.3</td>
</tr>
<tr>
<td>$u \rightarrow d \bar{(t)}$</td>
<td>5</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>$t$ quark</td>
<td>40</td>
<td>36</td>
<td>0</td>
</tr>
</tbody>
</table>

Table 2

The approximate cross sections and distributions are compared with the complete $2 \rightarrow 3$ reaction in Table 2. The agreement is satisfactory [indeed, not only for the average values but also for the differential distributions]. As a result, the $2 \rightarrow 3$ process $u + b \rightarrow W + d + t$ etc. can be used for $t$ production via $Wg$ fusion in shower Monte Carlos as Pythia and Herwig.

As will be studied in the experimental section in more detail, there is a large background to the final states in $Wg$ fusion from

$$pp \rightarrow W + jj + [j]$$

production, one of the jets being emitted at small transverse momentum. The corresponding cross sections have been evaluated in the matrix element approach as well as in parton shower Monte Carlos. The signal can only be isolated if the experimental resolution is so excellent that a peak in the $W + jet$ invariant mass distribution at the top mass can be seen.

Single top production in $Wg$ fusion as well as single-lepton events in $q\bar{q}/gg$ fusion to $t\bar{t}$ pairs can be exploited to measure the top quark mass independently of any theoretical model. The $W$ mass constraint $m_{W} = (p_{t} + p_{\nu})$ allows us to reconstruct the longitudinal $\nu$ momentum. There remains an ambiguity in the solution

$$p_{\nu, T} = \frac{1}{2p_{W}} \left[ p_{\nu T} \left( m_{W}^{2} + 2p_{\nu T}p_{T} + E_{\nu}m_{W}\sqrt{m_{W}^{2} + 4p_{\nu T}p_{T}} \right) \right] \tag{1.8}$$

However, in most cases ($\approx 65\%$) the lower value is physically realized as may be proved by comparing the solutions with Monte Carlo analyses.

2. Top Quark Fragmentation *

The hadronization process of top quarks is complicated because strong and weak mechanisms are intimately intertwined. The inclusive scale is set by the lifetime of the top quark [8], approximately given by

$$\tau_{t}^{-1} \approx 175 \text{ MeV} \times \left( \frac{m_{t}}{m_{W}} \right)^{3} \tag{2.1}$$

for sufficiently heavy top masses in the $SM$ decay $t \rightarrow b + W^{+}$. [Details will be discussed in the next section.] For $m_{t} \leq 100 \text{ GeV}^{-1}$ the lifetime $\tau_{t} \geq \Lambda^{-1}$ is long enough to form non-perturbatively mesonic $T_{t} \approx (0q)$ or baryonic $(tqg)$ bound states. The perturbative gluon radiation at early times is followed by a small non-perturbative deceleration of the $t$ hadron [14-17]. On the other hand, if $m_{t} \geq 100 \text{ GeV}$, the decay is so rapid that hadrons cannot build-up anymore. The top quark is decelerated only by the early radiation of hard, non-collinear gluons [8].

Perturbative gluon radiation off heavy quarks is described by leading order by the non-singlet Altarelli-Parisi equation [18]. The solution for the $t$ energy spectrum [scaled by the initial parton energy $E$]

$$< x >_{PQCD} = \left[ \frac{a_{s}(E)}{a_{s}(m_{t})} \right]_{\mu = \Lambda} \approx \frac{1}{x} \left[ \frac{a_{s}(E)}{a_{s}(m_{t})} \right]_{\mu = \Lambda} \tag{2.2}$$

reveals that only a small fraction of energy is radiated off for an initial energy $E$ above several $m_{t}$ in the $gg$ fusion process [$N_{t} = 5$]. For a top mass of 150 GeV and $\Lambda = 200$ MeV, the average $< x >_{PQCD}$ falls from 0.92 to 0.90 when the initial parton energy rises from 500 GeV to 1 TeV. The angular distribution $(\Theta)$ and the energy distribution $(\omega)$ of the radiated gluon is approximately given by

$$dF_{\omega} = \frac{4\alpha_{s}}{3\pi} \frac{G_{F}}{\sqrt{\omega}} \frac{d\Theta}{\Theta} d\omega \tag{2.3}$$

for a short-lived gluon radiation source accelerated to $\gamma = E/m_{t}$. The gluons accumulate on the surface of a cone with half-aperture $\Theta_{t} \sim \gamma^{-1}$ for a long-lived $t$ and $\gamma^{-1} \sqrt{\gamma} \gamma_{\lambda} / \omega$ if the particle decays quickly. The spectrum has a maximum at $\omega \sim \gamma_{\lambda} / \gamma$ if $\gamma_{\lambda}$ rises beyond the

*V. Khose and P. M. Zerwas
confinement energy $\Lambda$. Below the maximum, the fragmentation function can be approximated by the simple leading log form [19]

$$D(z)_{QCD} \propto \frac{1 + z^2}{(1-z)^{1+\eta}} \quad \text{with} \quad \eta = \frac{16}{33 - 2N_f} \log \left[ \frac{\alpha_s(E^2)}{\alpha_s(m_t^2)} \right]$$

(2.4)

This leading log form is improved by terms of order $\alpha_s \log(1-z)$ and $\alpha_s \log^2(1-z)$, generating a maximum in the energy distribution that is followed by a sharp fall-off when $z$ approaches 1 [20].

For small top quark masses $m_t < 100$ GeV the $t$-hadron energy is derived by convoluting the perturbative spectrum with the non-perturbative fragmentation function so that

$$<z >_t = <z >_{NP} <z >_{QCD}$$

(2.5)

The non-perturbative deceleration is small for heavy top quarks. Adopting the Peterson et al. form for the fragmentation function [17] we find

$$<z >_{NP} = 1 - \sqrt{E_t} \quad \text{with} \quad \sqrt{E_t} \approx 400 \text{ MeV}/m_t$$

(2.6)

The non-perturbative fragmentation function itself is strongly peaked near $z \approx 1$.

For top quark masses $m_t \geq 100$ GeV the weak decay process $t \rightarrow b + W^+$ intercepts the fragmentation process [8]. $t$ bound states do not form anymore, and small hadronization effects due to the beginning of string formation [as long as $m_t < (2\sqrt{m_t})^{1/2}$] are shrouded by perturbative gluon radiation. Also the perturbative gluon radiation is cut-off for lack of time at $\omega \geq \Gamma_t / \gamma_t$. The top fragmentation is therefore described by a narrow peak at $z \approx 1$, followed by the standard QCD continuum distribution.

After the decays of the top quarks in the $gg$ fusion process, $t\bar{t} \rightarrow b\bar{b}W^+W^-$, the bottom quarks fragment independently of each other. Their color flux lines are connected to the gluonic remnants of the perturbative $t$ and $\bar{t}$ fragmentation processes [21].

3. Top Quark Decay in the Standard Model

Based on the Standard Model, the lower limit of the top quark mass $m_t > 89$ GeV has been shifted to a level well above the real $W$ threshold for the decay mode $t \rightarrow b + W$. In the 3-family SM the CKM matrix element of the $(tb)$ coupling $V_{tb} = 0.9990 \pm 0.0004$ is well-known to be very close to unity [22]. As a result, the width of the top quark can be precisely predicted in the SM [8]:

$$\Gamma(t \rightarrow b + W^+) = \frac{G_F m_t^3}{8\pi^2} \left[ 1 - \frac{m_b^2}{m_t^2} \right] \left[ 1 + 2 \frac{m_W^2}{m_t^2} \right] \left( 1 + m_W^2/m_t^2 \right)$$

(3.1)

The width approaches quickly the asymptotic form

$$\Gamma(t \rightarrow b + W^+) \rightarrow 175 \text{ MeV} \times \left( \frac{m_t}{m_W} \right)^3$$

(3.2)

and the lifetime drops below $10^{-22}$ sec. This is illustrated in Fig.9. The final state consists of a $b$-quark jet plus a lepton pair or two other jets from the $W$ decay which, being a color singlet system, fragment independently of the environment.

With rising top quark mass the $W$ boson is emitted more and more frequently in a state of longitudinal polarization [23]. The ratio of longitudinal polarization to transverse polarization $\Gamma_L/\Gamma_T$ rises from 1/2 at the real $W$ threshold to $2m_t^2/m_W^2$ asymptotically. If the decaying top quark is polarized with a degree $P_t$, the charged $t^\pm$ lepton is emitted preferentially in the top spin direction with a probability $\frac{1}{2} \left[ 1 + P_t \cos \theta \right]$.

The energy spectra of the leptons, and correspondingly the invariant $(bl)$ mass, $\mu^2 = m_b^2 + l^2/m_t^2$, are affected by the chirality of the $(tb)$ charged weak current for moderate top masses. Assuming the charged lepton current to be of the well established $V-A$ form, the probability to observe a charged lepton of maximal energy $E_l = m_t/2$ vanishes for the conventional $V-A$ form of the $(tb)$ current while it is non-zero for a speculative $V+A$ coupling. Equivalently, the distribution of the invariant $(bl)$ mass is given by

$$\frac{1}{\Gamma} \frac{d\Gamma}{d\mu^2} \propto \begin{cases} \mu^2(1 - \mu^2) + \eta(1 - 2\mu^2) - \eta^2 & (V-A)_b \text{ coupling} \quad (SM) \\ \mu^2(1 - \mu^2) & (V+A)_b \text{ coupling} \quad (3.3) \end{cases}$$

where $\eta = m_W^2/m_t^2$ and $0 \leq \mu^2 \leq 1 - \eta$. A deviation of the $(tb)$ coupling from the $V-A$ structure of the SM would thus lead to a stiffening of the charged lepton spectrum.

Genuine electroweak corrections to the top quark decay are expected to be of order $G_F m_t^3/M^3$ while QED corrections, summed up and no energy and angular cuts applied, should be $O(\alpha/\pi)$. 

*P. M. Zerwas
QCD gluonic corrections, though being well under control, reduce the width nevertheless at a level of ~ 10%. Furthermore, they distort the charged lepton spectrum of the W decay [24]. The correction to the width can be cast into a simple form

$$\Gamma = \Gamma_0 \left[ 1 - \frac{2}{3} \frac{a_s}{\pi} f \right]$$

(3.4)

with, for $m_b = 0$,

$$f = \frac{F_1}{F_0} \quad F_0 = 2 (1 - \eta)^3 (1 + 2 \eta)$$

$$F_1 = F_0 \left[ \pi^2 + 2 S \eta (1 - \eta) - 2 S (1 - \eta) \right] + 4 \eta (1 - \eta - 2 \eta^2) \log \eta$$

$$+ 2 (1 - \eta)^3 (5 + 4 \eta) \log (1 - \eta) - (1 - \eta) (5 + 2 \eta - 6 \eta^2)$$

In the mass range $m_t \geq 100$ GeV, the QCD correction amounts to

$$\delta_{QCD} \approx (3 \text{ to } 4) a_s/\pi$$

(3.5)

The change of the lepton spectrum [described by a somewhat clumsy formula] may be taken from the literature directly [24].

4. Top Signals and Backgrounds

A. General features of signals

The most promising top signals via $t \to bW$ decays are [25,26]

$$gq/q\bar{q} \to t \bar{t} \to b\bar{b}WV \to b\bar{b}l\nu j$$

(4.1)

$$ug/\bar{u}g \to t \bar{t} \to b\bar{b}WV \to b\bar{b}l\nu l$$

(4.2)

$$ug \to b\bar{b} \to d\bar{b}W \to d\bar{b}l\nu$$

(4.3)

The essential ingredients in these signals are

i) Charged leptons with high $p_T$, usually isolated from accompanying hadrons.

ii) Neutrinos from $W \to l\nu$, giving high missing transverse energy $E_T$. In single-lepton cases, the transverse mass $m_T(l, E_T) = [2p_T(l) \cdot E_T]^{1/2}$ has a Jacobian peak near $m_T = m_W$, smeared by measurement errors and other contributions to $E_T$.

iii) b and \bar{b} jets from $W \to b\bar{b}$. They have typically very high $p_T$ with a Jacobian peak at $p_T(b) = (m_t^2 - m_W^2)/2m_t$, smeared by top transverse motion. For $m_t > 150$, this peak has $p_T(b) > 50$ [we use GeV units throughout]. The essential $b$-flavor can be tagged by a decay muon in the jet or by microvertex detection of displaced decay vertices.

iv) Jets from $W \to jj$, with high $p_T$ (Jacobian peak at $p_T \sim 1/2 m_W$, smeared by $W$ transverse motion); $p_T$ is typically less than b-jets, for $m_t > 150$. Jet pairs have $m(jj) \sim m_W$.

B. Backgrounds with similar features

i) The largest QCD backgrounds are from b and c quark pairs:

$$QCD \to b\bar{b} (or c\bar{c}) X \to \text{leptons} + E_T + \text{jets}.$$  

(4.4)

\*R. J. N. Phillips

For high-$p_T$ leptons, $b\bar{b}$ dominates over $c\bar{c}$. [High-$p_T$ production is similar, but $b$ has harder fragmentation and harder lepton decay spectrum.]

The first step is to demand high $p_T(l)$: for $t \to lX$ signals we find

$$p_T(l) \geq 30 : \quad \text{background/signal} \sim 40 (700) \text{ for } m_t = 150 (300)$$

$$p_T(l) \geq 60 : \quad \text{background/signal} \sim 5 \text{ (70) for } m_t = 150 (300)$$

(4.5)

Another step is to require isolation (e.g. sum of hadronic energy to be less than 8 GeV in a cone ($\Delta R^2 = (\Delta\eta)^2 + (\Delta\phi)^2 < 0.16$ about the lepton direction). This typically suppresses the background by a factor ~ 25 (50) for $p_T(l) \geq 30$ (60).

For single-lepton signals, we can require $m_T(l, E_T)$ to lie near $m_W$, e.g. $50 \leq m_T \leq 90$: This introduces a suppression factor ~ 20.

For dilepton signals, we can require that the two leptons are not back-to-back e.g. $30^\circ < \Delta\phi(l_1 l_2) < 150^\circ$, since this background is strongly peaked near $\Delta\phi = 0^\circ$ and $180^\circ$ while the $t\bar{t}$ signals are not.

This background contains two b-jets, but any high-$p_T$ leptons are correlated with them, so they cannot fully fake this aspect of the signals.

ii) Genuine $W$ production [27]:

$$pp \to W + jets \to l\nu + jets$$

(4.6)

This fakes the leptonic aspects of one-lepton signals [ $p_T(l), E_T$ and $m_T(l, E_T)$ ]. But the accompanying jets are mostly less hard than those from $t \to bW$ and $W \to jj$. Furthermore, they have no peak at $m(jj) \sim m_W$ and they rarely have $b$-flavor.

iii) Double-$W$ production [28]:

$$pp \to W + jets \to l_1 l_2 \nu \bar{\nu} + jets$$

(4.7)

This fakes the leptonic side of two-lepton signals, but not the jets [see (ii)]. $WW$ production by 4-parton scattering and event overlap are negligible compared to the direct process $q\bar{q} \to WW$.

iv) Drell-Yan electroweak dilepton production:

$$pp \to (\gamma^*, Z \rightarrow e^+ e^-, \mu^+ \mu^-) + jets$$

(4.8)

This gives isolated dileptons with a peak at $m(l_1 l_2) = m_Z$. The leptons always have the same flavor, $E_T$ is small (from measurement errors etc.); for jets see (ii).

v) $\tau$ pair production:

$$pp \to (\gamma^*, Z \rightarrow \tau^+ \tau^-) + jets$$

(4.9)

$\tau$ decays give isolated leptons or jets plus $E_T$. The leptons can have any flavor ($e$ or $\mu$) but are relatively soft; there is no $m(l_1 l_2)$ constraint; $E_T$ can be large. However, if we form the cluster transverse mass, from the cluster $c = l_1 + l_2$,

$$m_T^2 = (p_T^2(c) + m_T^2)^{1/2} + E_T^2 - (p_T^2(c) + E_T^2)$$

we get $m_T < m_Z$ [modulo measurement errors etc. contributing to $E_T$ and off-shell $Z$, $\gamma$ tail].
C. Specific signals

i) $\ell \ell \rightarrow 1$ lepton + ($n \leq 4$) jets.

The $b\bar{b}$ background requires lepton cuts. Requiring e.g., $p_T(l) > 50$ with $50 < m_T(l, E_T) < 90$ and lepton isolation probably suppresses this below the signal for $m_t < 300$ [see (B) above for numbers]. The $pp \rightarrow W + n$ jet background * can be reduced but not removed by savage jet cuts. For $n \leq 3$ this background process has been available already for some time. The $n = 4$ case has recently been evaluated [27]. The many subprocesses can be classified according to the number of participating quarks: 2, 4 or 6 quarks. The two-quark subprocesses contain many gluons. It then becomes profitable to use a technique recursive in the number of gluons. For the other processes helicity amplitudes are calculated using the Weyl-van der Waerden spinor calculus. For the evaluation of the $W + n$ jet cross section the MRSEB parton densities are used and $M_W$ is taken as the QCD scale. The following cuts on the transverse energy $E_T$ of the jets and the lepton, on the missing transverse energy, on the pseudorapidity $\eta$ of the jets and the lepton and on the separation $\Delta R$ between two jets or between a jet and a lepton were applied to both the signal and the background calculation:

\[
E_T^{\text{min}}(j) = E_T^{\text{min}}(l) = E_T^{\text{min}}(\text{miss}) = 50 \text{ GeV}
\]

\[
|\eta^{\text{max}}(j)| = |\eta^{\text{max}}(l)| = 3
\]

\[
\Delta R^{\text{min}}(j, j) = \Delta R^{\text{min}}(j, l) = 0.4
\]

The results are listed in Tables 3 and 4. Table 3 gives the total background process for $n = 1, \ldots, 4$ jets and its division into quark subprocesses. It should be noted that the relevance of the four-quark processes increases with $n$. The six-quark process is still negligible.

In Table 4 the signal as obtained from [29] is given for various $m_{top}$ values together with the background. If in the background only the events containing a $b\bar{b}$ are selected, the background effectively disappears.

\[
\begin{array}{cccc}
\hline
n & 2 \text{ quarks} & 4 \text{ quarks} & 6 \text{ quarks} & \text{total} \\
\hline
1 & 80 & - & - & 80 \\
2 & 45 & 7.0 & - & 52 \\
3 & 18 & 5.8 & - & 24 \\
4 & 6.1 & 2.4 & 0.13 & 8.6 \\
\hline
\end{array}
\]

We can also search for the top signal as a peak above background, e.g. in the invariant mass distribution of jets in the hemisphere opposite to the trigger lepton [30].

ii) Electroweak $u \bar{u} \rightarrow d \bar{b} \rightarrow$ isolated lepton + ($n \geq 3$) jets.

This channel becomes competitive for very heavy top, $m_t > 300$ as discussed in the previous section. The $b\bar{b}$ and $W +$ jets backgrounds are serious, as above. But here we have a new trick (see [13]); the $b$-jet from $t \rightarrow bW$ is expected to have the smallest rapidity among the high-$p_T$ jets. With $b$ identified and $\nu$ reconstructed (from $E_T$ and $m(\ell\nu)$, we can construct the invariant mass $m(b\ell\nu)$; the top signal can in principle be seen as a narrow peak above background in this mass distribution. Unfortunately this requires very accurate $E_T$ measurements.

iii) $t\bar{t} \rightarrow 2$ isolated leptons + ($n \geq 2$) jets.

A lepton $p_T$ cut brings the $b\bar{b}$ background in sight:

\[
p_T(l), E_T > 40 \quad : \quad b\bar{b} \quad \text{background/signal} \sim 5\; (50) \quad \text{for} \quad m_t = 150\; (300)
\]

(4.10)

$\Delta \phi(l, b)$ and isolation cuts in both leptons then suppress this background. The $W$ is then removed by $m_T(l, E_T) > m_Z$.

The $WW +$ jets background remains, being similar to the signal in the lepton sector. Requiring 2 hard jets removes it: e.g. $p_T(j) > 50$ with $p_T(l), E_T > 25$ gives

\[
\sigma (WWjj \rightarrow \ell \nu\nu jj) \sim 0.04 \; \text{pb} \tag{4.11}
\]

\[
\sigma (t\bar{t} \rightarrow \ell \nu\nu jj) \sim 9\; (1.5) \; \text{pb} \quad \text{for} \quad m_t = 150\; (300)
\]

This leaves a clean signal. Additional $b$-tagging is not necessary, but would make the signal extremely clean and help in reconstructing observed top events.

\footnote{F. Berends}
5. Non-Standard Model Decays

5.1 The Charged Higgs Decay Mode of the Top Quark

Among possible extensions of the Standard Model the minimal supersymmetric model MSSM has received much attention [31]. For this model to be anomaly-free, two Higgs doublets must exist, generating altogether 5 physical particles. Two of the particles are charged, \( H^\pm \), two scalars \( h, H \) and one pseudoscalar \( A \) \( \text{[CP]} = -1 \) are neutral. At the tree level, the mass spectrum is restricted by the relations

\[
m_{H^\pm}^2 = m_W^2 + m_A^2
\]

\[
m_{h, H}^2 = \frac{1}{2} \left( m_Z^2 + m_H^2 - \sqrt{(m_Z^2 + m_H^2)^2 - 2m_W^2m_H^2 \cos^22\beta} \right) \leq m_Z^2
\]  (5.1)

The angle \( \tan \beta = \frac{v_2}{v_1} \) is given by the ratio of the vacuum expectation values of the Higgs fields which provide the masses to up and down quarks, respectively. In the MSSM, \( v_2/v_1 \) is expected [32] to acquire a value between 1 and \( \sim m_t/m_b \). LEP analyses have set lower limits on the light neutral Higgs particle \( m_A > m_h > 32 \text{ GeV} \) (for a summary see Ref.[33]). Lower limits on \( m_h \) restrict the parameter range of \( |m_{H^\pm}, \tan \beta| \) to the areas shown in Fig. 10.

Fig. 10. Allowed areas of \( |m_{H^\pm}, \tan \beta| \) for given light scalar Higgs masses in the MSSM. [Below the dashed lines the decay mode \( H^\pm \to bW^\pm \) is kinematically forbidden.]

For a wide range of parameters the decay of top quarks to charged Higgs bosons \( t \to b + H^\pm \) is kinematically allowed and, as a result of the large quark masses involved, the magnitude of the branching ratio is noticeable [8,31,34,35]. A comparison of the width

\[
\Gamma (t \to b + H^\pm) = \frac{G_F m_t^3}{4\sqrt{2} \pi} \left( 1 + \frac{m_b^2 - m_f^2}{m_t^2} \right)^{\frac{1}{2}} \left[ \left( \frac{m_t}{m_b} \right)^2 \tan^2 \beta + \cot^2 \beta \right] + 4 \left( \frac{m_b}{m_t} \right)^3
\]  (5.2)

Fig. 11. Branching ratios for top decays to \( W^\pm \) and charged Higgs \( H^\pm \). The minimal value of \( \tan \beta \) in the MSSM depends on the light Higgs mass and can be read off the previous figure.

Fig. 12. Branching ratios of charged Higgs decays for \( m_{H^\pm} = 100 \text{ GeV} \) where the decay to \( W \) and light Higgs is kinematically forbidden. The lower limit for \( \tan \beta \) in the MSSM can be read off Fig. 10.

*P. M. Zerwas and J. Zunft
with the $SM$ width in eq. (3.1) shows that in the physically interesting range of $\tan \beta$ the Higgs decay does not occur. We look for the search strategies for top quark pairs at sufficiently high mass for those Higgs particles in the $SM$ that decay into a top quark pair and Higgs mass of 150 GeV.

Charged Higgs boson decay into heavy fermion pairs $H^+ \rightarrow t \bar{b}$ is large enough to be searched for charged Higgs particles in top quark mass of 150 GeV and a Higgs mass of 150 GeV. This is illustrated quantitatively in Fig. 11 for a top quark mass of 150 GeV and a Higgs mass of 150 GeV. 

They add up to a small value of $\tan \beta$ (less than 100 GeV). Excited particles for low $\tan \beta$ and large $m_H$ are very narrow in the parameter range relevant for charged Higgs boson decay. The branching ratio is

$$\Gamma (H^+ \rightarrow t \bar{b}) = \frac{G_F}{4 \pi \sqrt{2}} \frac{m_H^2}{2} \left( \frac{m_t^2}{m_H^2} \right)^2 \left( 1 - \frac{m_t^2}{m_H^2} \right)^2 \tan \beta \cos \beta$$

(3.1)

and $\tan \beta = \sin \beta / \cos \beta$ as well as $W_A$ and $H_W$ are allowed by the Higgs boson decay.

For large $\tan \beta$, the decay rate is nearly unchanged, providing an excellent signature for charged Higgs boson decay. The branching ratio is

$$\Gamma (H^+ \rightarrow t \bar{b}) = \frac{G_F}{4 \pi \sqrt{2}} \frac{m_H^2}{2} \left( \frac{m_t^2}{m_H^2} \right)^2 \left( 1 - \frac{m_t^2}{m_H^2} \right)^2 \tan \beta \cos \beta$$

(3.2)

and $\tan \beta = \sin \beta / \cos \beta$ as well as $W_A$ and $H_W$ are allowed by the Higgs boson decay.

5.2 Supersymmetric Extension - Stop

Alternatively, if the LSP is a scalar (as in the Standard Model), the dominant decay mode will be $H^+ \rightarrow t \bar{b}$, which gives rise to a soft lepton energy spectrum and to a more pronounced transverse mass distribution. In this case there will be no hard primary lepton in the final state.

Addendum. Because of low rates and large backgrounds it is very difficult to search for direct $H^+ \rightarrow t \bar{b}$ pair production through the electroweak Drell-Yan mechanism in $pp$ collisions. Besides the top decay to $b \bar{b}$, which is associated with the forward production of top and charged Higgs in the range of $1 \to 10$ GeV, the Higgs particle in the $H^+$ channel [37].

In supersymmetric (SUSY) extensions of the Standard Model, the two scalar partners of the left- and right-handed fermions give rise to a $2 \times 2$ mass matrix, including the $b$ quark, this mixing is immaterial. For the heavy quark, however, it could become significant.
To summarize, the typical signatures of SUSY t-decays are (i) larger missing energy and (ii) softer charged lepton spectra than in the SM t-decay. At present there are too many unknown SUSY parameters for conclusive analyses and unambiguous predictions. Therefore, the best strategy to search for these SUSY signals in top-decays would be to look for access events which are not compatible with the precise expectations for the SM decay $t \rightarrow bW$. On the other hand, the observation of a t quark conforming exactly to the SM expectations would serve to rule out light stops and/or neutralinos! A more detailed experimental feasibility study is presented in the SUSY Section of these proceedings.

6. Non-SM Quarks

6.1 Fourth generation quarks

If a fourth generation of leptons and quarks $(\nu_t, L, (t', b')$ exists [48], the quark masses must exceed 89 GeV to escape the CDF bound [49] and the lepton masses must be of order $m_Z/2$ or greater to conform with the LEP data. The mass splittings are constrained by electroweak radiative corrections [50]

$$(m_t - m_{t'})^2 + \frac{1}{3} (m_{t'} - m_b)^2 + m_{\nu_t}^2 \leq 200 \text{ GeV}^2$$

(5.6)

QCD production of $b'b'$ and $t't'$ pairs is similar to $t\bar{t}$; electroweak $t'\bar{b}'$ production is suppressed relative to the $tb$ case, however, since both quarks are heavy. If $b'$ were lighter than $t$ [51], it would be produced more copiously and it would be discovered first. Its probable decay mode $b' \rightarrow b\gamma$ gives signals very similar to top; however, the lepton and neutrino distributions are different compared to $t \rightarrow bW$ (essentially interchanged), and this may eventually be discernible [52]. If $m_t > m_{t'} + m_b$, $t'$ can decay by $t' \rightarrow b\gamma t \rightarrow \gamma W W$, so that $t'\bar{b}'$ production gives distinctive four-W final states. Different mass orderings and mass splittings lead to a variety of scenarios that we do not pursue in detail here; there are many similarities [30] to the standard $t\bar{t}$ scenario, but also significant differences do exist.

6.2 Exotic quarks

New quarks may be present below 1 TeV and they could be produced in this mass range at large hadron colliders as shown in the Summary. If they do not belong to standard families, they are called exotic. The simplest and most attractive possibility is a new vector-like quark of charge $-\frac{2}{3}$ in each family. The left and right-handed components of vector-like fermions have the same transformation properties under $SU(2)_L \times U(1)_Y$. Such quarks [53] are present in many extensions of the Standard Model, for instance in extended gauge models, in models of spontaneous $CP$ violation, etc. Present data constrain their masses little because their mixing with the standard quarks is naturally small.

We discuss the possibility to detect such quarks at the LHC, assuming that the matter content of the theory is that of the minimal Standard Model plus this new fermion.

For a large range of the parameter space a new vector-like D quark will have similar signatures as the top quark. Its dominant decay is into a $W$ plus a charge $\frac{2}{3}$ quark (typically 50\%) which looks like a top decay if we cannot determine the flavor of the final quark. However, $D$ can also decay to $Z$ or $H$ (typically 25\% of each of them) plus a light quark. Therefore, although $D$ should be searched for in the same final states as the top, its character will be determined by observing its neutral decay modes. The CDF lower limit $m_D > 89$ GeV applies approximately to $m_D$ too. However, whereas the partial analysis of neutral current data and LEP results (including radiative corrections) indicate $m_t < 200$ GeV, $m_D$ has no upper bound.

![Fig.13. $\sigma$ (pb) as a function of the $D$ quark mass (in TeV). All transverse momenta ($l, j, m$) missing are required to be $> 50$ GeV, whereas the top mass is equal to 200 GeV.](image)

We have calculated the $pp \rightarrow DD$ cross sections plus the decays into the most interesting channels, (a) $W^+W^-q\bar{q} \rightarrow e\nu\bar{\nu}j\bar{j}$, (b) $e\nu\bar{\nu}j\bar{j}$ and (c) $ZWq\bar{q} \rightarrow e\nu\bar{\nu}j\bar{j}$. The analysis is carried out at the parton level and the normalization is chosen as $\sigma(pp \rightarrow DD) = \sigma(pp \rightarrow t\bar{t}) = 0.58$nb for $m_D = m_t = 200$ GeV; this is in the middle of the QCD predictions. In all cases we require $p_T^{\text{missing}} > 50$ GeV. We assume the $DD$ decay fractions to be (a) $(\frac{1}{2})^2 \times 2 \times \frac{1}{2} = 0.67$, (b) $(\frac{1}{2})^2 \times 2 \times (\frac{1}{2})^2$ and (c) $2 \times \frac{1}{2} \times \frac{1}{2} \times 0.034 \times \frac{1}{2}$ where the $D$ decay fractions into $W$ and $Z$ are equal to $\frac{1}{2}$ and $\frac{1}{2}$, respectively. We implicitly assume that the Higgs is lighter than the new quark, and the D decay fraction into $H$ is $\frac{1}{2}$.

In channels (a) and (b), the $DD$ signal is like $t\bar{t}$ but reduced by a factor $\frac{1}{2}$. The two cases are therefore difficult to distinguish if $m_t < 200$ GeV as expected. The horizontal lines marked $t\bar{t}$ in Fig.13 show the $t\bar{t}$ backgrounds for $m_t = 200$ GeV. [For other backgrounds and references see the $t\bar{t}$ discussions.]
In channel (c) the $t \bar{t}$ background enters via semi leptonic $b$-decays; if we impose an additional mass cut $|m(e^{+}e^{-}) - m_Z| < 10$ GeV, the cross section is about 0.008 pb as shown in Fig.13. A further isolation cut (not shown) could reduce the $t \bar{t}$ background by a factor of 50, approaching the irreducible $pp \rightarrow ZWjj$ background which is estimated to be a few times $10^{-5}$ pb.

We conclude that the $t \bar{t}$ background (let alone other backgrounds) renders channels (a) and (b) unpromising, but that a $DD$ signal should be detectable in the NC channel (c) up to $m_D \sim 400 - 500$ GeV for $f \mathcal{L} = 10^{3}$ pb$^{-1}$, and eventually up to 700 GeV for a ten times larger integrated luminosity.

REFERENCES

    V. Fadin, V. Khoze and T. Sjöstrand, CERN-TH 5687/90.
    Y. Dokshitzer et al., in preparation.
    and Wesley 1990) and references quoted therein.
[32] See the contribution by F. Zwirner [SUSY Group] to these proceedings.
[33] F. Dydak, Proceedings, XXV. Intern. Conference on High Energy Physics,
    Singapore 1990.

[37] J. Zunft, contribution to this Workshop; 
A. C. Bawa et al., contribution to this Workshop, based on Ref.[35].

H. E. Haber and G. L. Kane, Phys. Rep. C117 (1985) 75;  
L. J. Hall, ibid., 197;  


[47] H. E. Haber and G. L. Kane, in Ref.[38].


D. Haidt, DESY 89-073;  
P. Langacker, UPR-0435 T (1990);  


V. Barger, N. G. Deshpande and J. F. Gunion et al., Proc. of 1986 Snowmass Workshop;  