CERN TH.5601/89

**Semi-leptonic Flavour Changing Neutral Current Decays
of The Fourth Generation b' -quark**

WEI-SHU HOU

*Paul Scherrer Institute
CH-5232 Villigen PSI, Switzerland*

and

ROBIN G. STUART

*Theory Division, CERN
CH-1211 Geneva 23, Switzerland*

ABSTRACT

We study the decay modes $b' \rightarrow b\ell^+\ell^-$, $b' \rightarrow b\nu\bar{\nu}$ for $m_{b'} \lesssim M_Z$ including virtual γ , Z^0 , and box diagram contributions, the latter complicating the computation considerably. Due to the effect of the Z^0 propagator and the finite Z^0 width, these decay rates are substantial. The box diagram is subleading but also not negligible. As a consequence, the lepton pair mass distribution is rather flat and high mass $\ell^+\ell^-$ or $\nu\bar{\nu}$ pairs are much more abundant than suggested from phase space. It is found that $b' \rightarrow bZ^*$ (loosely denoting the sum of $b' \rightarrow b\ell^+\ell^-$, $b' \rightarrow b\nu\bar{\nu}$ and $b' \rightarrow bq\bar{q}$) is of similar order as the $b' \rightarrow bg$ decay mode, and is likely to dominate in case $m_{b'}$ is very heavy. With three light neutrinos, the decay $b' \rightarrow b\nu\bar{\nu}$ is of similar order as $b' \rightarrow b\gamma$, however, the $b' \rightarrow b\ell^+\ell^-$ modes are relatively suppressed. The flavour changing neutral current b' decay scenario is thus shown to be quite rich. A comprehensive, but theoretical, discussion is given regarding b' discovery potentials at present and future colliders.

CERN-TH.5601/89

December 1989

1. Introduction

In previous papers [1, 2] we pointed out that if the fourth generation charge $-\frac{1}{3}$ quark, b' , is lighter than the top quark then, under reasonable assumptions about the structure of the extended quark mixing matrix, flavour changing neutral current (FCNC) processes could be a significant and even the dominant decay mode. It must be emphasized that this rather unusual situation arises entirely within the Standard Model extended to four generations. With the top quark appearing to be very heavy, the only additional assumption, that b' is lighter than top, becomes increasingly plausible. The search for the fourth generation b' -quark decaying via FCNC must be continued until such time as the top quark has been discovered and the mass region below completely explored.

The immediate consequence of this is clear. The usual experimental signature of "isolated lepton plus jets", appropriate for charged current decays is no longer applicable. It was shown that one can then look for peaking in the mass spectrum of 4-jet events from the decay $b' \rightarrow b\gamma$ or for events containing an isolated photon + three jets when one of the b' -quarks undergoes the decay, $b' \rightarrow b\gamma$. A number of experimental groups have made searches based on our suggestions [3-8]. The MARK II collaboration has looked for the above signatures and set a combined limit of 42.0 GeV on the b' mass at the 95% confidence level [3]. Currently the best experimental mass limit, regardless of the dominant decay modes, is 45.2 GeV [8]. This value is obtained by the analysis of event shape distributions, particularly acoplanarity, rather than a direct search for 4-jet or 3-jet+photon events; a luxury that is only possible in the e^+e^- environment and with the large data sample of LEP compared to SLC. Such searches rely on the relatively high sphericity of decays of the slowly moving heavy quark. In machines where the event shape distributions can be used to detect heavy fermions, the signatures mentioned above are complementary and independent measurements that would facilitate the study of any observed variations.

To date no b' searches that assume FCNC dominance have been made at hadron colliders. Analyses have looked only for the charged current (CC) decay modes [9] and are therefore unable to discover the b' or place useful limits on its mass in the case where FCNC decays dominate CC decays by a substantial amount. The

CDF collaboration at the TEVATRON has the potential to explore the highest mass region, but with their new data sample, UA2 could soon discover or exclude the b' up to a mass of around 60-70 GeV.

Recent measurements of the Z^0 resonance width at SLC and LEP [10] seem to favour the existence of three light neutrino species rather than four. Despite bold proclamations in the popular press this measurement does not unambiguously exclude a fourth generation. It must be remembered that the Z^0 line shape is at best an indirect measurement of the neutrino number, N_ν . If new physics beyond the Standard Model is present it could, in principle, modify the lineshape in such a way as to mask the presence of the additional neutrinos. Of course such new physics would eventually affect other precision measurements via radiative corrections, but the number of independent measurements that probe the Standard Model to this level of sensitivity is at present very limited. Further, the lineshape itself is not directly sensitive to neutrinos with a mass greater than about $M_Z/2$. It would be naive in the extreme to abandon heavy fermion searches based on such indirect evidence. In fact the observation of the b' from its FCNC decays could be the first signal of the fourth generation.

In Refs. 1, 2 a detailed study of the various FCNC decay modes $b' \rightarrow b\gamma$, $b\gamma$ and bZ was carried out. The calculation is complicated by the relatively large number of independent but non-negligible masses. Essential use was made of the algebraic manipulation package LERG-I [11, 12] written in REDUCE to handle the reduction of expressions to a minimal form, perform stringent algebraic checks and finally write FORTRAN code for numerical evaluation. The conclusions can be summarized as follows: Below the Z^0 production threshold $b' \rightarrow b\gamma$ is the dominant FCNC decay mode with $b' \rightarrow b\gamma$ having the remarkably large relative branching fraction of up to around 20%. Above Z^0 threshold the mode $b' \rightarrow bZ$ quickly becomes by far the dominant mode.

In the present paper we report on an extension of our original calculation to FCNC decay modes involving leptonic final states $b' \rightarrow b\ell^+\ell^-$ and $b' \rightarrow b\nu\bar{\nu}$. At first sight, these modes would be expected to be significant only for $m_{b'} \lesssim M_Z$ and may provide a clean signal of the b' . Surprisingly, it turns out that the partial decay rate induced by virtual γ , Z^0 and box diagrams, which we shall generically

call $b' \rightarrow bZ^*$, are found to be comparable to the on-shell $b' \rightarrow bg$ mode. Thus, our earlier results for the rates $b' \rightarrow bg$ and $b' \rightarrow b\gamma$ are diminished by roughly a factor of two or more. The $b' \rightarrow bZ^*$ modes themselves lead, however, to characteristic signatures that still may be useful for the purposes of analysis, after a discovery has been made, especially in the hadronic environment.

The layout of the paper is as follows: In Section 2 we discuss the computation of the box diagram. FCNC vertex corrections have been considered already in Ref. 2. Section 3 contains explicit expressions for the products of matrix elements relevant to our calculation. Section 4 deals briefly with possible checks that can be made. For the box diagram these are rather limited compared to those available for the FCNC vertices. In Section 5 we present our results and in Section 6 we give intuitive arguments to explain the relative strength of $b' \rightarrow bZ^*$ processes. Experimental implications are discussed. The appendix explains the notation for 4-point functions used in the calculation of the box diagram.

2. Calculation of the Box Diagram

The Feynman diagrams contributing to the processes $b' \rightarrow b\nu\bar{\nu}$ and $b\ell^+\ell^-$ are shown in Fig. 1a&b. The sum of all FCNC $b'bZ$ and $b'b\gamma$ vertex diagrams are represented by the shaded blobs. They have been studied in detail in Ref. 2. The box diagrams, introducing as they do, the need to handle 4-point functions add considerably to the complexity of the calculation both algebraically and numerically. They cannot simply be dropped as they are necessary to yield gauge invariant S -matrix elements and their contribution is also found to be numerically significant.

Consider the box diagram contributing to the process $b' \rightarrow b\nu\bar{\nu}$. As in our previous studies, we will work in the Euclidean metric with time-like momenta being negative. Conversion to the popular Minkowskian metric of Bjorken and Drell [13] is straightforward and is discussed in detail in Ref. 2. We define $p_{b'}$ to be the incoming 4-momentum of the b' -quark and p_b , p_ν , $p_{\bar{\nu}}$ to be the outgoing momenta of the b -quark, the neutrino and anti-neutrino respectively. The Mandelstam variable s is the invariant mass of the $\nu\bar{\nu}$ pair while u is the invariant mass of the outgoing fermions b , ν . The mass of the internal heavy quark will be generically labelled m_i .

One has the useful identities [14],

$$[\gamma_\mu \gamma_\sigma \gamma_\nu \gamma_\tau \gamma_\rho \gamma_\lambda \gamma_\kappa]_1 [\gamma_\nu \gamma_\rho \gamma_\mu \gamma_\tau \gamma_\lambda \gamma_\kappa]_2 = 4 [\gamma_\rho \gamma_\tau \gamma_\lambda \gamma_\kappa]_1 [\gamma_\sigma \gamma_\nu \gamma_\mu \gamma_\tau]_2, \quad (2.1a)$$

$$[\gamma_\mu \gamma_\sigma \gamma_\nu \gamma_\tau \gamma_\rho \gamma_\lambda \gamma_\kappa]_1 [\gamma_\nu \gamma_\rho \gamma_\mu \gamma_\tau \gamma_\lambda \gamma_\kappa]_2 = 4 \delta_{\rho\sigma} [\gamma_\alpha \gamma_\tau \gamma_\lambda \gamma_\kappa]_1 [\gamma_\alpha \gamma_\nu \gamma_\mu \gamma_\tau]_2, \quad (2.1b)$$

$$[\gamma_\mu \gamma_\sigma \gamma_\nu \gamma_\tau \gamma_\rho \gamma_\lambda \gamma_\kappa]_1 [\gamma_\mu \gamma_\rho \gamma_\nu \gamma_\tau \gamma_\lambda \gamma_\kappa]_2 = 4 \delta_{\rho\sigma} [\gamma_\alpha \gamma_\tau \gamma_\lambda \gamma_\kappa]_1 [\gamma_\alpha \gamma_\nu \gamma_\mu \gamma_\tau]_2, \quad (2.1c)$$

$$[\gamma_\mu \gamma_\sigma \gamma_\nu \gamma_\tau \gamma_\rho \gamma_\lambda \gamma_\kappa]_1 [\gamma_\mu \gamma_\rho \gamma_\nu \gamma_\tau \gamma_\lambda \gamma_\kappa]_2 = 4 [\gamma_\rho \gamma_\tau \gamma_\lambda \gamma_\kappa]_1 [\gamma_\sigma \gamma_\nu \gamma_\mu \gamma_\tau]_2, \quad (2.1d)$$

where $\gamma_{L,R}$ is the left/right hand helicity projection operator and the square brackets, $[\]_1$, $[\]_2$, indicate that the enclosed quantities act on different fermion lines. Note that the in the identities (2.1a-d) terms proportional to $(n-4)$ have been dropped, where n is the dimensionality of space-time that occurs naturally in dimensional regularization. Since the box diagrams considered here are UV finite this can be safely done.

With the aid of eq. (2.1a) the contribution of the neutrino box diagram is easily shown to be,

$$\mathcal{M}_{\text{box}}^{b' \rightarrow b\nu\bar{\nu}} = \frac{g^4}{16\pi^2} [\gamma_\alpha \gamma_L]_\nu [\gamma_\alpha \gamma_L]_b F_4, \quad (2.2)$$

with,

$$F_4 = \left\{ -\frac{1}{2}(u - m_b^2)D_{11} - \frac{1}{2}(u + m_b^2)D_{12} + \frac{1}{2}(u + m_b^2)D_{13} \right. \\ \left. - m_b^2 D_{22} - m_b^2 D_{23} - (u - m_b^2)D_{24} \right. \\ \left. + (s + u - m_b^2)D_{25} - (s - m_b^2 - m_b^2)D_{26} + 4D_{27} \right\}, \quad (2.2a)$$

where $[\]_\nu$, $[\]_b$ act on neutrino and quark states respectively. The 4-point form factors D_{ij} are defined in the appendix following the conventions of Passarino and Veltman [15]. Their arguments are, $D(0, -m_b^2, -m_b^2, 0, -u, -s; 0, M_W^2, M_W^2)$. In practice we set $m_b = 0$ when evaluating form factors which dramatically shortens the final results. m_b is kept non-zero where it appears in the calculation of products of matrix elements and in phase space.

For the process $b' \rightarrow b\ell^+\ell^-$ define s , as before, to be the invariant mass of the outgoing lepton pair and t to be the invariant mass of the outgoing b -quark and anti-lepton, ℓ^+ . If the outgoing lepton has 4-momentum, p_{ℓ^-} and that of the anti-lepton

is p_l^+ then the box diagram now has the more general form,

$$\mathcal{M}_{\text{box}}^{b' \rightarrow b e^+ e^-} = -\frac{g^4}{16\pi^2} \left\{ F_1 [\bar{p}_b \gamma_L] d [\bar{p}_{e^-} \gamma_L] b + F_2 [\bar{p}_b \gamma_L] d [\bar{p}_b \gamma_L] b \right. \\ \left. + F_3 [\bar{p}_b \gamma_L] d [\bar{p}_b \gamma_L] b + F_4 [\gamma \alpha \gamma_L] d [\gamma \alpha \gamma_L] b \right\}, \quad (2.3)$$

in which the quantities in square brackets $\{ |e, | b$ act respectively on the lepton and quark states and,

$$\begin{aligned} F_1 &= -D_{24} + D_{25} - D_{11}, & (2.3a) \\ F_2 &= D_{22} - D_{26} + D_{12}, & (2.3b) \\ F_3 &= D_{23} - D_{26} + D_{13}, & (2.3c) \\ F_4 &= D_{27}. & (2.3d) \end{aligned}$$

The arguments of the form factors, D_{ij} , are in this case,

$$D(0, -m_b^2, -m_b^2, 0, -t, -s; 0, M_W^2, m_1^2, M_W^2).$$

The expressions (2.2) and (2.3) are now in a form suitable for input into the package LERG-I that algebraically reduces the expressions to a few scalar integrals, B_0 , C_0 and D_0 . These forms have the advantage that they are unique and that many large cancellations, which are potential sources of numerical instability, have been removed at the outset. The output obtained from LERG-I can be obtained directly in the form of FORTRAN code thereby eliminating the danger of transcription errors.

The calculation described so far has been performed in the Euclidean metric; conversion to Minkowskian metric is simply a matter of the substitution $D_{27} \rightarrow -D_{27}$.

A detailed study on the evaluation the scalar integrals, B_0 , C_0 and D_0 can be found in Ref. 16. For the first two, relatively simple general expressions are known and have been programmed. The general form of D_0 is extremely complicated and for its evaluation we use the relation,

$$D_0(p_1^2, p_2^2, p_3^2, p_4^2, p_5^2, p_6^2; 0, m_2^2, m_3^2, m_4^2) = A_2 A_3 A_4 C_0(k_1^2, k_2^2, k_3^2, k_5^2, M_2^2, M_3^2, M_4^2), \quad (2.4)$$

in which,

$$A_2 = \frac{1}{p_1^2 + m_2^2}, \quad A_3 = \frac{1}{p_5^2 + m_3^2}, \quad A_5 = \frac{1}{p_4^2 + m_4^2}, \\ M_i^2 = A_i^2 m_i^2,$$

$$\begin{aligned} k_1^2 &= A_2^2 p_1^2 - A_2 A_3 (p_1^2 - p_2^2 + p_3^2) + A_3^2 p_5^2, \\ k_2^2 &= A_4^2 p_4^2 - A_3 A_4 (-p_3^2 + p_4^2 + p_6^2) + A_3^2 p_5^2, \\ k_5^2 &= A_2^2 p_1^2 - A_2 A_4 (p_1^2 + p_4^2 - p_6^2) + A_4^2 p_4^2, \end{aligned}$$

With D_0 defined as,

$$D_0(p_1^2, p_2^2, p_3^2, p_4^2, p_5^2, p_6^2; m_1^2, m_2^2, m_3^2, m_4^2) = \frac{1}{\int d^n q} \frac{1}{[q^2 + m_1^2][(q + p_1 + p_2)^2 + m_2^2][(q + p_1 + p_2 + p_3)^2 + m_3^2]}, \quad (2.5)$$

$$p_1 + p_2 + p_3 + p_4 = 0, \quad p_5 = p_1 + p_2, \quad p_6 = p_2 + p_3,$$

the relation (2.4) is obtained by the substitution $q_\mu = r_\mu/r^2$ in the integral (2.5). The relation (2.4) must however be used with considerable care and is not an exact identity in all regions of parameter space. The reason is that the above substitution does not faithfully preserve the $i\epsilon$ prescription. That is to say, the infinitesimal imaginary parts of the arguments M_2^2 , M_3^2 and M_4^2 of C_0 in (2.4) are not necessarily all negative. It is also not valid for complex masses since it will generate complex momenta k_1 , k_2 , k_5 . However, below the physical thresholds of the C_0 and for at least one external non-space-like momentum, the $i\epsilon$ terms only affect the imaginary parts of the form factor. These will be zero overall when the D_0 in eq. (2.4) is evaluated below its thresholds. Thus it is an identity provided that D_0 is evaluated below threshold and the imaginary part of the right hand side of eq. (2.4) is discarded. In addition, at least one of the effective momenta, k_1 , k_2 , k_5 must be either light-like or time-like. Fortunately this is just the region of interest for the particular calculation described here, $m_b \leq m_i$. The threshold $m_b > 2M_W$ is also far outside the region of study here. Although we have not proved that it is generally true, we do check that the constraint on k_1 , k_2 and k_5 is empirically satisfied during numerical integrations.

3. Matrix Element Products

The box diagram increases the number of possible forms that can occur in the products of matrix elements. In this section, general expressions are given for the possible matrix element product that appear. The b -quark mass is kept because in the evaluation of phase space integrals numerical instabilities can often be tamed by keeping small masses that act as cutoffs.

The general single particle exchange diagram containing a FCNC vertex contributing to $b' \rightarrow b \ell^+ \ell^-$ (see Fig. 1) can be represented for massless leptons as,

$$\mathcal{M}_{\text{vertex}} = [V_\ell]_\ell [V_q]_b P, \quad (3.1)$$

where

$$\begin{aligned} V_q &\equiv \gamma_\mu (F_L^q \gamma_L + F_R^q \gamma_R) + i \sigma_{\mu\nu} q_\nu (F_L^q \gamma_L + F_R^q \gamma_R), \\ V_\ell &\equiv \gamma_\mu (F_L^\ell \gamma_L + F_R^\ell \gamma_R), \end{aligned}$$

and,

$$q = p_{b'} - p_b = p_{\ell^+} + p_{\ell^-},$$

P is the effective internal propagator, and the subscript ℓ here denotes either charged or neutral lepton current. From now on we revert to the Minkowskian metric. With $\mathcal{M}_{\text{box}}$ as in (2.3), but dropping the factor $-g^4/16\pi^2$, the required products of matrix elements become,

$$\begin{aligned} |\mathcal{M}_{\text{box}}|^2 = & \left\{ |F_1|^2 (s + u - m_b^2)(u - m_b^2) + |F_2|^2 (-s + m_b^2 + m_{b'}^2) m_{b'}^2 \right. \\ & + |F_3|^2 (-s + m_b^2 + m_{b'}^2) m_b^2 + 2 \operatorname{Re}(F_1 F_2^*) (u - m_b^2) m_{b'}^2 \\ & + 2 \operatorname{Re}(F_1 F_3^*) (s + u - m_b^2) m_b^2 + 4 \operatorname{Re}(F_1 F_4^*) (u - m_b^2) \\ & + 4 \operatorname{Re}(F_2 F_3^*) m_b^2 m_{b'}^2 + 4 \operatorname{Re}(F_2 F_4^*) m_{b'}^2 + 4 \operatorname{Re}(F_3 F_4^*) m_b^2 \Big\} \\ & \times (-su - u^2 + um_b^2 + um_{b'}^2 - m_b^2 m_{b'}^2) \\ & - 4 |F_4|^2 (u - m_b^2)(u - m_{b'}^2) \end{aligned} \quad (3.2)$$

$$\begin{aligned} \mathcal{M}_{\text{vertex}} \mathcal{M}_{\text{box}}^* = & 2 F_L^\ell \left\{ F_L^q F_1^* (u - m_b^2) + F_L^q F_2^* m_{b'}^2 + F_L^q F_3^* m_b^2 \right. \\ & + F_R^q F_2^* m_b m_{b'} + F_R^q F_3^* m_b m_{b'} + F_{TL}^q F_1^* s m_b \\ & - F_{TR}^q F_2^* s m_{b'} - F_{TL}^q F_3^* s m_b \Big\} \\ & \times (-su - u^2 + um_b^2 + um_{b'}^2 - m_b^2 m_{b'}^2) \\ & - 4 F_L^\ell F_4^* \left\{ F_L^q (u - m_b^2)(u - m_{b'}^2) + F_R^q s m_b m_{b'} \right. \\ & \left. + F_{TL}^q s m_b (u - m_{b'}^2) + F_{TR}^q s m_{b'} (u - m_b^2) \right\} \end{aligned} \quad (3.3)$$

$$\begin{aligned} |\mathcal{M}_{\text{vertex}}|^2 = & -4 \left\{ |F_L^q F_L^\ell|^2 + |F_R^q F_R^\ell|^2 (u - m_b^2)(u - m_{b'}^2) \right. \\ & + |F_R^q F_L^\ell|^2 + |F_L^q F_R^\ell|^2 (s + u - m_b^2)(s + u - m_{b'}^2) \\ & + |F_L^\ell|^2 + |F_R^\ell|^2 (F_L^q F_R^{q*} + F_L^q F_R^q) s m_b m_{b'} \\ & + |F_L^\ell|^2 (F_L^q F_{TL}^{q*} + F_L^q F_{TL}^q) s m_b (u - m_{b'}^2) \\ & + |F_R^\ell|^2 (F_R^q F_{TR}^{q*} + F_R^q F_{TR}^q) s m_b (u - m_{b'}^2) \\ & + |F_L^\ell|^2 (F_L^q F_{TL}^{q*} + F_L^q F_{TL}^q) s m_{b'} (u - m_b^2) \\ & + |F_R^\ell|^2 (F_R^q F_{TR}^{q*} + F_R^q F_{TR}^q) s m_{b'} (u - m_b^2) \\ & - |F_L^\ell|^2 (F_R^q F_{TL}^{q*} + F_R^q F_{TL}^q) s m_{b'} (s + u - m_{b'}^2) \\ & - |F_R^\ell|^2 (F_L^q F_{TR}^{q*} + F_L^q F_{TR}^q) s m_{b'} (s + u - m_b^2) \\ & - |F_L^\ell|^2 (F_R^q F_{TR}^{q*} + F_R^q F_{TR}^q) s m_b (s + u - m_b^2) \\ & - |F_R^\ell|^2 (F_L^q F_{TL}^{q*} + F_L^q F_{TL}^q) s m_b (s + u - m_{b'}^2) \\ & - |F_L^\ell F_L^q|^2 + |F_R^\ell F_R^q|^2 (s + u - m_{b'}^2)(u - m_b^2) \\ & - |F_L^\ell F_R^q|^2 + |F_R^\ell F_L^q|^2 (s + u - m_b^2)(u - m_{b'}^2) \\ & + |F_L^\ell|^2 + |F_R^\ell|^2 (F_L^q F_{TL}^{q*} + F_L^q F_{TL}^q) s^2 m_b m_{b'} \Big\} \end{aligned} \quad (3.4)$$

The more general $\mathcal{M}_{\text{vertex}}^*$ can easily be obtained from eq. (3.4).

It is a straight forward matter to include the coupling factors that have been dropped in the foregoing analysis and include the factor $(8\pi m_{b'})^{-3}$ to give the 3-

body differential decay rates in s and u or t . To calculate total decay rates the Monte Carlo integration routine VEGAS was used.

4. Checks

Ref. 2 described in detail, stringent checks that were applied to the FCNC vertices $b'bZ$, $b'bg$ and $b'by$. The most severe coming from current conservation of the vector parts of these vertices. Unfortunately no such powerful check of this nature is available for the box diagram. One could calculate both box and vertex diagrams in a general R_ξ gauge and look for cancellation of the gauge parameter between vertex and box diagrams but it would add enormously to the complexity of the calculation. Comparison with known low energy results is possible but does not provide a severe test because most previous calculations have been done with zero mass external fermions, in which case only the D_{27} form factor remains for both the neutrino and charged lepton boxes. We have, nevertheless, confirmed that our result reduces to that of $b \rightarrow s\ell^+\ell^-$, $b \rightarrow s\nu\bar{\nu}$ [17] for zero mass external fermions. Note also, however, that the expressions that are used as input to LERG-I for reduction to scalar integrals are quite simple and can be obtained with very little hand manipulation.

The reduction of the form factors F_1 , F_2 , F_3 and F_4 to expressions involving only the scalar integrals B_0 , C_0 and D_0 generates denominators that are powers of the product $s \cdot t \cdot u$, for $m_b = 0$, and is small in certain phase space limits. The near vanishing of this denominator is compensated for by cancellations between various scalar integrals. Numerical stability of our results when evaluated near kinematic limits is a good indication that the expressions are correct and that B_0 , C_0 and D_0 are all being evaluated correctly numerically.

Although it will not be presented here, we have also checked that the decay rate for $b' \rightarrow bZ^*$ with a constant width of 2.5 GeV included in the internal Z^0 propagator, maps smoothly onto that of $b' \rightarrow bZ$ calculated in Ref. 2 for $m_Y > m_b + M_Z$. For further discussion of this point see Ref. 18.

5. Results

The results of the calculation are shown in Fig. 2-6. Throughout the following discussion M_Z is set to 91.1 GeV, $M_W = 80$ GeV and $\sin^2 \theta_W = 1 - M_W^2/M_Z^2$. Fig. 2a shows the partial decay rate for the process $b' \rightarrow c$ or $t + e\bar{\nu}$ for the ratio of KM matrix elements $V_{cb'}/V_{tb'} = 1, 0.1, 0.01, 0.001$ (from top to bottom), and for $m_t = 80$ and 120 GeV. An overall factor of $|V_{tb'}|^2$ has been removed. The sharp upturn in all curves at $m_Y \sim 80$ GeV is due to the physical W threshold. Below $b' \rightarrow t$ threshold, curves with $m_t = 80$ GeV and $m_t = 120$ GeV are identical. The lowest curve ($V_{cb'}/V_{tb'} = 0.001$) shows a bifurcation near this threshold, however, above this point FCNC decays are less likely to dominate for $m_t = 80$ GeV. The other curves also bifurcate but not within the frame of the plot.

For comparison Fig. 2b shows the FCNC decay rate for $b' \rightarrow b\ell^+\ell^-$ without the KM factor $|V_{tb'}^*V_{tb'}|^2$. The dashed curves correspond to $m_Y = 300$ GeV and the solid ones to $m_Y = 200$ GeV. In both pairs the upper and lower curves are calculated for $m_t = 80$ GeV and 120 GeV respectively. The upper is truncated at $m_Y = m_t$ above which our hypothesis of FCNC dominance begins to lose validity and there would be computational difficulties for evaluating D_0 from eq. (2.4). Fig. 2c shows a corresponding set of curves for the decay $b' \rightarrow b\nu\bar{\nu}$, summing over three light neutrino species.

A similar set of graphs are shown in Fig. 3 but with $m_t = 120$ and 150 GeV, and for higher m_Y values. In Fig. 3a the curve for $|V_{cb'}/V_{tb'}| = 1$ is outside the frame of the graph. For $|V_{cb'}/V_{tb'}| = 0.001$ the $m_t = 120$ GeV curve branches sharply upward from the $m_t = 150$ GeV curve when the decay $b' \rightarrow te\bar{\nu}$ becomes kinematically allowed. The beginnings of bifurcation can be seen in the $|V_{cb'}/V_{tb'}| = 0.01$ curve above. In Fig. 3b and c dashed curves correspond to $m_Y = 350$ GeV and the solid curve to $m_Y = 250$ GeV. The sharp rise near the lower end of the m_Y scale is due to the turning on of the on-shell Z^0 boson (see Fig. 1), much like the case of the turn on of on-shell W bosons in Fig. 2a. A finite width of 2.5 GeV has been included in the Z^0 propagator. Note that we have again summed over three light neutrino species in Fig. 2c. Still, the total $b' \rightarrow b\nu\bar{\nu}$ rate is roughly 10-11 times larger than the single $b' \rightarrow be^+e^-$ channel. The decay rates to $\mu\bar{\mu}$ and $\tau\bar{\tau}$ are comparable with the $e\bar{e}$ mode within the given experimental errors. In fact, the $e\bar{e}$ mode is expected

to be slightly larger due to the photonic F_2 contribution. After removing a factor of 3, and another factor of roughly 2 to account for the difference between the ν and e Z -charge, the two rates still differ by a factor of about 1.8. This is due to a combination of reasons. The box diagram tends to enhance the $b' \rightarrow b\nu\bar{\nu}$ rate by roughly 40% or more, while suppressing the $\ell^+\ell^-$ mode by roughly 10% or more. The remainder is mostly due to the extra photonic contribution for $b' \rightarrow b\ell^+\ell^-$, which tends to interfere destructively with the rest. The upshot is a sizeable $b' \rightarrow b\nu\bar{\nu}$ mode, whereas the mode with better signature, $b' \rightarrow b\ell^+\ell^-$, is rather suppressed.

Fig. 4 shows the relative rates for the FCNC decays a) $b' \rightarrow bg$, b) $b' \rightarrow b\gamma$ and c) $b' \rightarrow bg^*$, where the notation is the same as in Fig. 2b and c, and as before the results are plotted without the factor $|V_{tb}^*V_{t\ell'}|^2$. One can compare these earlier results with Fig. 2. Note that in Fig. 4c only the $g^* \rightarrow q\bar{q}$ contribution due to the “charge radius” F_1 form factor have been included. This is because a low-energy theorem seems to indicate the absence of contributions from F_1 to the process $g^* \rightarrow g\gamma$ [19]. Although, of course, the low-energy theorem itself has no reason to be applicable here, it does suggest that the F_1 contribution to $b' \rightarrow bgg$ may be quite subtle. We argue below that the virtual gluon effect is not important at all.

With three light neutrinos, the result given in Fig. 2c for $b' \rightarrow b\nu\bar{\nu}$ is comparable to the photonic $b' \rightarrow b\gamma$, and may even be larger in case ℓ' is rather heavy. As stated in Ref. 2, and can be readily seen by comparing Fig. 5a and b, the photonic $b' \rightarrow b\gamma$ is more than 10% of the $b' \rightarrow bg$ rate. Thus, the rate for $b' \rightarrow b\nu\bar{\nu}$ being of comparable size means that the *total* effect of virtual Z^0 and box diagram contribution cannot be ignored. To this effect, one needs to calculate $b' \rightarrow bq\bar{q}$ coming from virtual γ , Z^0 and box diagrams, with q being the usual five species of light quarks. The contribution of the box diagram could be obtained from the results given in this paper for all but the decay $b' \rightarrow bbb$, where the presence of an additional heavy internal quark complicates matters. Here only estimates are presented as the uncertainties already involved, in particular in $m_{b'}$, are rather forgiving. Taking the relative strength of $Z \rightarrow q\bar{q}$ vs. $Z \rightarrow \nu\bar{\nu}$ decay as a guide, we find that the total contribution of $b' \rightarrow b\ell^+\ell^-$, $b\nu\bar{\nu}$, $bq\bar{q}$ (which we denote from now on $b' \rightarrow bZ^*$), to be about 4.5 times as large as the result for $b' \rightarrow b\nu\bar{\nu}$ in Fig. 2c. Thus, it is seen that $b' \rightarrow bZ^*$ is comparable to $b' \rightarrow bg$, and even dominating when $m_{b'}$ is large. An immediate consequence of this is that 6-jet events arising from the production of $b'\bar{b}'$ pairs will

be comparable to 4-jets events. A discussion of implications for search strategies will be given later. It is clear, however, that the photon, Z^0 , plus box diagram contribution to $b' \rightarrow bq\bar{q}$ will completely dominate over the $b' \rightarrow bg^*$ contribution of Fig. 4c.

As further handles for the $b' \rightarrow b\ell^+\ell^-$ and $b' \rightarrow b\nu\bar{\nu}$ modes, we show in Fig. 5 and Fig. 6 the distribution of $d\Gamma(b' \rightarrow b\ell^+\ell^-)/dy$ and $d\Gamma(b' \rightarrow b\nu\bar{\nu})/dy$ vs. $y = s/m_{\ell\ell'}^2$ for $m_t = 100$ GeV and $m_{b'} = 200$ GeV, and $s \equiv q^2$ is the invariant mass of the outgoing lepton pair. Once again the KM factor $|V_{tb}^*V_{t\ell'}|^2$ is not included. $m_{b'}$ is 85 GeV for the upper curve, 65 GeV for the middle and 45 GeV for the lower. Note in particular here the singularity at $q^2 = 0$ for $b' \rightarrow b\ell^+\ell^-$ coming from the photon exchange graph that is absent in the case of final state neutrinos. The singularity is logarithmic and occurs due to the “dipole moment” form factor, F_2 in the notation of Ref. 2, of the photon FCNC vertex. This is not surprising since the same form factor leads to the on-shell $b' \rightarrow b\gamma$ process, so the virtual photon contribution would show a logarithmic mass singularity in the differential rate. In practice, although it is cut off at the phase space limit, $q^2 = 4m_e^2$, it would still lead to an enormously sharp peak in Fig. 5. We have only plotted down to y values of 10^{-3} . Another noticeable feature in Figs. 5 and 6 is their flatness in q^2 . There is an approximate combined suppression factor of $(1-y)^2$ coming from matrix element squared and phase space, which can be readily seen in the analogous case of $b \rightarrow s\ell^+\ell^-$ (see, e.g. Fig. 1 of Ref. [20]). In the present case, this is largely compensated for by the Z^0 propagator, and the fact that the Z^0 is unstable and so has a finite width. Thus, given that $m_{b'}$ has been taken to be comparable but smaller than M_Z , the propagator effect favours the large q^2 region whereas the matrix element squared and phase space effect favours not so large q^2 values. The combined effect is the flatness of the q^2 dependence up to rather high lepton pair masses, the more so the heavier the b' . This effect should help cut against background when searching for these decay modes.

6. Discussion and Conclusion

Let us stress again the physical reason why the apparently second order (electro)weak processes $b' \rightarrow b\ell^+\ell^-$, $b\nu\bar{\nu}$, $bq\bar{q}$ sum up to be comparable to the first order weak, first order strong process $b' \rightarrow b\gamma$, and may even be dominant over the tree level, first order weak, process $b' \rightarrow cW^*$. The former processes are dominated by the effective FCNC mediated by the virtual Z^0 boson. The $Zb\bar{b}'$ vertex form factor is very strongly dependent on the heavy quark masses m_t , $m_{t'}$. However, since the dependence is close to quadratic, (i.e. approximately $m_{t'}^2 - m_t^2$, so long that m_t is not too close to $m_{t'}$) the $m_{t'}$ contribution dominates and one is not very sensitive to m_t . The heavy $m_{t'}$ roughly compensates for the extra " G_F " suppression introduced by the Z^0 propagator. In contrast to the corresponding process, $b \rightarrow s\ell^+\ell^-$, $\nu\bar{\nu}$ case, for the range of b' masses under consideration, the proximity of $m_{t'}$ to M_Z has weakened this G_F factor considerably already. Thus, the combined effect is that the apparently second order weak process is effectively close to first order. In the scenario that we present, the CC $b' \rightarrow cW^*$ is likely to be KM suppressed, while $b' \rightarrow bZ^*$ is loop suppressed. From this perspective, the fact that $b' \rightarrow bZ^*$ wins over $b' \rightarrow cW^*$ when $V_{cb'} \lesssim 0.01$ becomes quite understandable, since the loop diagrams are subject to suppression factor of roughly $1/(16\pi^2)$. As for the gluon induced FCNC, the main difference is gauge invariance. The fact that the gluon couples only to conserved currents leads to rather severe form factor suppression; the $m_{t'}$ dependence above is replaced at most by a factor of $m_{t'}$. The dependence of the gluon F_2 (dipole) form factor on heavy t and t' is even weaker than $\log(m_{t'}^2/m_t^2)$, which amounts to further suppression given that m_t is not dramatically different from $m_{t'}$. Thus, although it is only first order in weak and strong couplings, is at same loop level as $b' \rightarrow bZ^*$, and enjoys a 2-body phase space rather than 3-body, $b' \rightarrow b\gamma$ turns out to be only comparable to $b' \rightarrow bZ^*$. The latter even has the potential to become the dominant FCNC decay in case $m_{t'}$ is very heavy.

An additional feature of the FCNC decay modes of the b' quark has thus been uncovered. It is of interest to contemplate the impact on search strategies at various colliders. As mentioned in the Introduction, experiments at SLC and LEP have recently set limits on b' almost up to the kinematic $M_Z/2$. For LEP experiments it was sufficient to just use general event shape analysis, due to the high statistics

enjoyed at LEP. It is interesting to note that this utilizes the dominance of $b' \rightarrow b\gamma$ in FCNC decays. Allowing the gluon to fragment, at α_s order, which does not amount to a huge suppression, it would allow the event to look more spherical and non-planar, even if one is close to the kinematical limit where the b' quarks carry little momentum. The observation that the $b' \rightarrow bZ^*$ modes are comparable or even dominant over $b' \rightarrow b\gamma$ would further enhance the power of the event shape analysis. As far as event shape is concerned, both CC and FCNC decay modes would proceed via 3-body processes, rather than the 2-body decays $b' \rightarrow b\gamma$ and $b' \rightarrow b\ell$, at least more than half of the time. Thus, regardless of the relative CC vs. FCNC branching ratio, one can still safely use a 3-body decay phase space in the event shape analysis without substantial loss. Hence, at LEP II energies, as the search continues, event shape analysis should suffer from no drawbacks.

To fully understand what kind of heavy quark one has produced, in case a signal in the event shape analysis is seen, one would need a rather sophisticated and detailed study of the decay products in spherical or non-planar events. The richness of the FCNC decay modes would thus serve to pinpoint the production of a b' quark, since current indications from CDF seem to imply that there is only a marginal possibility for the top quark to be discovered even at LEP II energies. A detailed study of all these decay modes would give us much information on the heavy b' , t' quarks, in particular their masses and mutual mixing. It is unfortunate that the smallness of $b' \rightarrow b\ell^+\ell^-$, being roughly $1/10$ of $b' \rightarrow b\gamma$ and $b' \rightarrow b\nu\bar{\nu}$ per charged lepton species, probably would not serve as a "discovery" mode, but rather as a "confirmation" mode. Even the $b' \rightarrow b\nu\bar{\nu}$ channel, although comparable if not larger than $b' \rightarrow b\gamma$, would have to be studied further with Monte Carlo simulations to see whether the large missing mass, together with three jets, can serve as a useful signature. Note, however, the FCNC decay dominance scenario still does not lead to a large fraction of "isolated $e/\mu + \text{jets}$ " events. In any case, our earlier conclusions regarding $b' \rightarrow b\gamma$ and $b' \rightarrow b\gamma$, although softened by $b' \rightarrow bZ^* \gtrsim b' \rightarrow b\gamma$, still serve as good signatures. Furthermore, the $b' \rightarrow bZ^*$ mode may lead to a sizeable enhancement of 6- or 5- jet vs. 4- jet or less events.

For hadronic machines, there is no such luxury as the utility of event shape analysis in the search for heavy quark production. In this case, one would have to utilize the FCNC decay signatures themselves. At the moment, only charge current

(isolated e/μ plus jets) signatures have been studied, concurrent with the search for the top. This is useful only when the CC decay mode has a substantial decay branching ratio (BR). As the charge current decay BR is lowered, at some level one would lose sensitivity. However, for instance, UA2 has demonstrated rather good capabilities for direct photon studies and if the direct photon event sample is relatively clean, one can look for the accompanying multi-jets (as opposed to the single back-to-back jet from Compton scattering) and its associated signatures. Note that our new result that $b' \rightarrow bZ^* \gtrsim b' \rightarrow bg$ would lead to $\gamma + 4$ -jets being comparable, if not larger than, $\gamma + 3$ -jets, which should serve as an interesting handle. The QCD background should have the opposite feature. Another possible handle is to look for heavy flavour tagging by muons inside jets that come from the b -quark decay. Given that UA1 has made good use of such in the search for $B-\bar{B}$ mixing, this approach should also be pursued. With regard to the TEVATRON, the most difficult region is that of low $m_{\ell\ell}$ mass because of lower p_T photons coming from b' decays. Given that SLC/LEP has set a limit of $m_{\ell\ell} \gtrsim M_Z/2$, this should help CDF to concentrate on the higher mass b' region. In this respect, the TEVATRON has great potential to explore the highest mass region presently available, perhaps even beyond LEP II energies. Aside from the $b' \rightarrow b\gamma$ signature, it remains to be seen if any of the $b' \rightarrow b\ell^+\ell^-$, $b' \rightarrow b\nu\bar{\nu}$, $b' \rightarrow bq\bar{q}$ modes can be of use in a hadronic environment.

In summary, we have studied the processes $b' \rightarrow b\ell^+\ell^-$, $b' \rightarrow b\nu\bar{\nu}$ for $b' \lesssim M_Z$ by a complete calculation. Contrary to naive expectations, these decays are quite substantial compared to the $b' \rightarrow b\gamma$, $b' \rightarrow bg$ modes suggested by us earlier. Using the results, we estimate $b' \rightarrow bq\bar{q}$ and find that $b' \rightarrow bZ^*$, which is the sum of this and the two other modes above, is of same order as $b' \rightarrow bg$, and has the potential to dominate in case ℓ' is rather heavy. Thus, the complete scenario for FCNC decay modes for $m_{b'} \lesssim M_Z$ is: $b' \rightarrow bZ^*$ is as important as $b' \rightarrow bg$, $b' \rightarrow b\nu\bar{\nu}$ is as important as $b' \rightarrow b\gamma$, while $b' \rightarrow b\ell^+\ell^-$ is approximately $1/10$ of the latter two for each charged lepton species, e, μ, τ , assuming that the fourth leptons δ and ν_δ are rather heavy. A possibly useful signature for the $b' \rightarrow b\ell^+\ell^-$, $\nu\bar{\nu}$ modes is the relative flatness of the lepton pair mass spectrum. Thus, one may perhaps look for the $b' \rightarrow b\nu\bar{\nu}$ mode by looking for events with missing mass of the order slightly less than M_Z with no isolated charged leptons, and 3-4 accompanying jets.

Event shape analysis would be good probes for the production of heavy flavours in an e^+e^- environment, whether they decay via FCNC or not. The FCNC decay modes are not only rich, but they may dominate the CC decay modes if the latter is strongly suppressed by small quark mixing angles. Even with the new SLC/LEP limits, there remains a vast domain in $m_{b'}$ below m_t for potential discovery. Of course, once $b' \gtrsim M_Z$ the on-shell $b' \rightarrow bZ$ process would begin to dominate the other FCNC decay modes, and the search strategy should be concentrated on Z^0 signals.

Acknowledgements

We thank D. Froidevaux and P. Igo-Kemenes for discussions.

Appendix A.

The scalar two-point function $B_0(p^2, m_1^2, m_2^2)$ and three-point function $C_0(p_1^2, p_2^2, p_3^2, m_1^2, m_2^2, m_3^2)$ are defined in Ref. 2 following the conventions of Passarino and Veltman [15]. In the present work we must also consider the scalar and related four-point functions defined in the Euclidean metric by,

$$D_0, D_\mu; D_{\mu\nu}(p_1, p_2, p_3; m_1, m_2, m_3, m_4) = \frac{\int \frac{d^n q}{i\pi^2} \frac{1; q_\mu; q_\mu q_\nu}{[q^2 + m_1^2][(q + p_1)^2 + m_2^2][(q + p_1 + p_2)^2 + m_3^2][(q + p_1 + p_2 + p_3)^2 + m_4^2]} \quad (A.1)$$

p_1, p_2, p_3 are external incoming 4-momenta and defining,

$$p_1 + p_2 + p_3 + p_4 = 0, \quad p_5 = p_1 + p_2, \quad p_6 = p_2 + p_3,$$

the required form factors $D_{ij}(p_1^2, p_2^2, p_3^2, p_4^2, p_5^2, p_6^2; m_1^2, m_2^2, m_3^2, m_4^2)$ are given by,

$$\begin{aligned} D_\mu &= (p_1)_\mu D_{11} + (p_2)_\mu D_{12} + (p_3)_\mu D_{13} \\ D_{\mu\nu} &= (p_1)_\mu (p_1)_\nu D_{21} + (p_2)_\mu (p_2)_\nu D_{22} + (p_3)_\mu (p_3)_\nu D_{23} + \{p_1 p_2\}_{\mu\nu} D_{24} \\ &\quad + \{p_1 p_3\}_{\mu\nu} D_{25} + \{p_2 p_3\}_{\mu\nu} D_{26} + \delta_{\mu\nu} D_{27} \end{aligned}$$

The braces $\{ \}$ here denote the symmetric product,

$$\{p_i p_j\}_{\mu\nu} = (p_i)_\mu (p_j)_\nu + (p_i)_\nu (p_j)_\mu$$

REFERENCES

- [1] W. S. Hou and R. G. Stuart, *Phys. Rev. Lett.* **62** (1989) 617.
- [2] W. S. Hou and R. G. Stuart, *Nucl. Phys. B* **320** (1989) 277.
- [3] G. S. Abrams *et al.* (MARK II collaboration), SLAC-PUB-5106, Oct. 1989
- [4] S. C. Eno, Ph.D. Thesis, Rochester University, (1989).
- [5] K. Abe *et al.* (VENUS collaboration), *Phys. Rev. Lett.* **63** (1989) 1776.
- [6] S. Eno *et al.* (AMY collaboration), *Phys. Rev. Lett.* **63** (1989) 1910.
- [7] I. Adachi *et al.* (TOPAZ collaboration), *Phys. Lett.* **229 B** (1989) 427.
- [8] M. Z. Akrawy *et al.* (OPAL collaboration), CERN EP 89.154.
- [9] D. Froidevaux, (UA2 collaboration), private communication
- [10] G. S. Abrams *et al.* (MARK II collaboration), *Phys. Rev. Lett.* **63** (1989) 724; B. Adeva *et al.* (L3 collaboration), *Phys. Lett.* **231 B** (1989) 509; D. Decamp *et al.* (ALEPH collaboration), *Phys. Lett.* **231 B** (1989) 519; M. Z. Akrawy *et al.* (OPAL collaboration), *Phys. Lett.* **231 B** (1989) 531; P. Aarnio *et al.* (DELPHI collaboration), *Phys. Lett.* **231 B** (1989) 539.
- [11] R. G. Stuart, *Comp. Phys. Comm.* **48** (1987) 367.
- [12] R. G. Stuart and A. G3ngora-T., preprint CINVESTAV/FIS-05/89 (to appear in *Comp. Phys. Comm.*)
- [13] J. D. Bjorken and S. D. Drell, *Relativistic Quantum Mechanics*, (McGraw-Hill New York, 1965).
- [14] A. Sirlin, *Nucl. Phys. B* **192** (1981) 93.
- [15] G. Passarino and M. Veltman, *Nucl. Phys. B* **160** (1979) 151.
- [16] G. 't Hooft and M. Veltman, *Nucl. Phys. B* **153** (1979) 365.
- [17] W. S. Hou, R. S. Willey and A. Soni, *Phys. Rev. Lett.* **58** (1987) 1608.

- [18] I. I. Bigi, Yu L. Dokshitser, V. Khoze, J. H. Kühn, and P. Zerwas, *Phys. Lett.* **181 B** (1986) 157.
- [19] J. Liu and Y.-P. Yao, U. Michigan preprint UM-TH-89-19, Oct. 1989.
- [20] G. W. S. Hou, *Int. J. Mod. Phys. A* **2** (1987) 1245.

FIGURE CAPTIONS

1. Feynman diagrams contributing to the flavour changing neutral current decays a) $b' \rightarrow b\nu\bar{\nu}$ and b) $b' \rightarrow b\ell^+\ell^-$. The shaded blobs represent the sum of FCNC vertex diagrams.
2. Partial decay rates of the b' -quark in GeV as a function of $m_{b'}$. a) $b' \rightarrow c$ or $t + e\bar{\nu}$ for the ratio of KM matrix elements $V_{cb'}/V_{tb'} = 1, 0.1, 0.01, 0.001$ (from top to bottom) and top quark masses 80 and 120 GeV. b) $b' \rightarrow b\ell^+\ell^-$; c) $b' \rightarrow b\nu\bar{\nu}$. In b) and c) the KM factor $|V_{tb'}^*V_{tb'}|^2 \approx |V_{tb'}^*V_{tb'}|^2$ has been removed. The dashed curves are for $m_t = 300$ GeV and the solid curves are for $m_t = 200$ GeV. In both sets the upper curve corresponds to $m_t = 80$ GeV and the lower to $m_t = 120$ GeV.
3. As in Fig. 2 except that $m_t = 120$ GeV and 150 GeV. In a) the curve for $V_{cb'}/V_{tb'} = 1$ lies outside the frame of the plot. In b) and c) the dashed curves are for $m_t = 350$ GeV and the solid curves are for $m_t = 250$ GeV. In both sets the upper curve corresponds to $m_t = 80$ GeV and the lower to $m_t = 120$ GeV.
4. Partial decay widths of the b' -quark in GeV, as a function of $m_{b'}$ with KM factors removed, a) $b' \rightarrow b\gamma$; b) $b' \rightarrow b\gamma$; c) $b' \rightarrow b\gamma$ excluding box diagram contributions.
5. Distribution of $d\Gamma(b' \rightarrow b\ell^+\ell^-)/dy$ as a function of $y = q^2/m_{b'}^2$ with KM factors removed. $q^2 \equiv s$ is the invariant mass of the lepton pair.
6. Distribution of $d\Gamma(b' \rightarrow b\nu\bar{\nu})/dy$ as a function of $y = q^2/m_{b'}^2$ with KM factors removed. $q^2 \equiv s$ is the invariant mass of the neutrino pair.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6