Proposing “$b$-Parity” - a New Approximate Quantum Number in Inclusive $b$-jet Production - as an Efficient Probe of New Flavor Physics

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Abstract

We consider the inclusive reaction $\ell^+\ell^- \to nb + X$ ($n =$ number of $b$-jets) in lepton colliders for which we propose a useful approximately conserved quantum number $b_P = (-1)^n$ that we call $b$-Parity ($b_P$). We make the observation that the Standard Model (SM) is essentially $b_P$-even since SM $b_P$-violating signals are necessarily CKM suppressed. In contrast new flavor physics can produce $b_P = -1$ signals whose only significant SM background is due to $b$-jet misidentification. Thus, we show that $b$-jet counting, which relies solely on $b$-tagging, becomes a very simple and sensitive probe of new flavor physics (i.e., of $b_P$-violation).

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The Standard Model (SM), despite its enormous success [1], is believed to be the low-energy limit of a more fundamental theory whose nature will be probed by the next generation of colliders [2]. New physics effects have been studied in a variety of processes using model independent approaches [3], as well as within specific models. All such investigations aim at providing a clear unambiguous signal for non SM physics effects which is also easily extracted from the experimental data. In this letter we propose one such signal which is obtained through simple $b$-jet counting. The approach is best suited to lepton colliders but may be extended to hadron colliders, e.g., it applies to the Fermilab Tevatron $p\bar{p}$ collider to the extent that the sea $b$-quark content of the protons can be ignored.

We will consider the inclusive multiple $b$-jet production in $\ell^+\ell^-$ collisions

$$\ell^+\ell^- \rightarrow nb + X,$$

where $n$ denotes the number of $b$ and $\bar{b}$-jets in the final state. For this type of reactions we introduce a useful approximate symmetry which we call $b$-Parity ($b_P$), defined as

$$b_P = (-1)^n.$$

In the limit where the quark mixing CKM matrix $V$ [4] satisfies $V_{3j} = V_{j3} = 0$ for $j \neq 3$, all SM processes are $b_P$-even since in this case the third generation quarks do not mix with the others, and this leads to the conservation of the corresponding flavor number. Given the fast top decay, only $b$-quarks and $t$-quark decay products are observed experimentally. Since $t \rightarrow bW$ with branching ratio $\sim 1$, the experimentally observed flavor number is in fact carried only by the $b$-quarks. Therefore, the measured quantum number reduces to the net number of detected $b$-quarks; we find it convenient to use instead the derived quantity $b_P$. 

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SM processes which violate this conserved number necessarily involve the charged current interactions \( \propto V_{ij} W^\mu \bar{u}_i \gamma_\mu (1 - \gamma_5) d_j \), where \( u_i \) and \( d_j \) denote up and down quarks of generations \( i \) and \( j \) respectively (with \( i \) or \( j = 3 \)); these are the only SM vertices that violate \( b_P \) (i.e., \( b_P = -1 \)).

In particular, reactions involving transitions of a \( b \)-quark to a \( c \) or \( u \)-quarks are proportional to \( |V_{cb}| \sim \lambda^2 \) and \( |V_{ub}| \sim \lambda^3 \) respectively, where \( \lambda \sim 0.22 \) is the Wolfenstein parameter \([5]\). The existing experimental numbers for these quantities are \( |V_{cb}| = 0.0397 \pm 0.002 \) and \( |V_{ub}/V_{cb}| = 0.08 \pm 0.02 \) \([6]\). Thus, all SM \( b_P = -1 \) processes are suppressed by these small off diagonal CKM elements \( |V_{cb}|^2 \) or \( |V_{ub}|^2 \); it follows that the SM is essentially \( b_P \)-even. As a consequence the irreducible SM background to \( b_P = -1 \) processes induced by new flavor physics beyond the SM, is strongly suppressed.

Note that in contrast with other observables, \( b_P \) relies solely on \( b \)-tagging not on the particular structure of a given final state, nor does it require the reconstruction of any other particle but the \( b \). Thus the main obstacle in this use of \( b_P \) is the reducible SM background where an odd number of \( b \)-jets are missed in an event with an even number of \( b \)-jets. This results from having a \( b \)-tagging efficiency \( (\epsilon_b) \) below 1. This type of background disappears as \( \epsilon_b \rightarrow 1 \), but even at \( \epsilon_b = 0.7 \) still produces a significant number of (miss-identified) events. The \( b \)-tagging efficiency is therefore the only relevant experimental parameter for this probe of new physics.

Consider now the inclusive \( b \) and \( \bar{b} \)-jets production process in (1). Since our method does not require the detection of the charge of the \( b \), \( n \) represents the number of \( b + \bar{b} \)-quarks in the final state. As mentioned above, \( b_P = 1 \) in the SM, since \( n \) is necessarily even. Thus, to estimate the background generated from missing an odd number of \( b \)-jets in a SM event we define the probability \( P^b_k \) for detecting \( k \) \( b \)-jets out of the \( n \) taggable \( b \)-jets in the final state

\[
P^b_k = \frac{n!}{k!(n-k)!} \epsilon_b^k (1 - \epsilon_b)^{n-k}.
\]
Although our method applies to any $b_P = -1$ processes, for simplicity, in what follows we focus only on $b_P = -1$ processes with $n = 1$ or $n = 3$ $b$-jets in the final state which provide the largest signal. To be specific we will consider reactions in $e^+e^-$ colliders, but the method is clearly extendable to muon colliders.

The leading SM $b_P = 1$ processes which generate a background to a $b_P = -1$ signal with $n = 1$ (i.e., to $e^+e^- \rightarrow b$-jet + $X$) are the following $2 \rightarrow 2$ reactions which give 2 or 4 $b$-jets in the final state

$$
\sigma_{bb} = \sigma(e^+e^- \rightarrow \bar{b}b), \quad \sigma_{tt} = \sigma(e^+e^- \rightarrow t\bar{t}), \\
\sigma_{ZZ} = \sigma(e^+e^- \rightarrow ZZ), \quad \sigma_{ZH} = \sigma(e^+e^- \rightarrow ZH). \quad (4)
$$

These cross-section are given for example in [7]. The $t\bar{t}$ final state produces 2 $b$-jets via the SM decay $t \rightarrow bW^+$ and its conjugate. The $ZZ$ and $ZH$ final states can give rise to 2 or 4 $b$-jets depending on whether one or two of these bosons decay to a $b\bar{b}$ pair.

Note also that there are no two-body processes in the SM with 6 $b$-jets in the final state. Thus, $\sigma_{ZZ}$ and $\sigma_{ZH}$ generate the only important $b_P = 1$ SM background to $b_P = -1$ processes with $n = 3$. We neglect $2 \rightarrow 3$ SM processes (which can also generate fake $b_P = -1$ signals) since these reactions are suppressed by small phase space factors. We note, however, that the SM $2 \rightarrow 3$ reactions will generate the dominating $b_P = 1$ (reducible) background to $n = 5$ $b_P = -1$ processes. The only SM irreducible background for $n = 1$ and $n = 3$ processes with $b_P = -1$ in $2 \rightarrow 2$ reactions comes from $e^+e^- \rightarrow W^+W^- \rightarrow b + X$ or $\bar{b} + X$ and $e^+e^- \rightarrow t\bar{t} \rightarrow b\bar{b}W^+W^- \rightarrow b\bar{b} + X$ or $b\bar{b} + X$; as noted before these processes involve a CKM-suppressed $W$-boson decays and the corresponding cross sections are unobservably small.

Using (3) the dominating background for $b_P = -1$ processes with $n = 1$ is given by

$$
\sigma_B^{(1)} = P_1^2 [\sigma_{bb} + \sigma_{tt} + 2B_Z\sigma_{ZZ} + (B_Z + B_H)\sigma_{ZH}] \\
+ P_1^4B_Z [B_Z\sigma_{ZZ} + B_H\sigma_{ZH}], \quad (5)
$$
where $B_Z$ and $B_H$ are the branching ratios for a $Z$-boson and a SM Higgs-boson to decay to a $b\bar{b}$ pair, respectively. Similarly, the only significant background for $b_P = -1$ processes with $n = 3$ is

$$
\sigma_B^{(3)} = P_3^4 B_Z (B_Z \sigma_{ZZ} + B_H \sigma_{ZH}) .
$$

In Figs. 1(a), (b) we plot the background cross-sections $\sigma_B^{(n)}$ for $n = 1, 3$ respectively, as a function of the $b$-tagging efficiency $\epsilon_b$ and for different c.m. collider energies. These cross-sections depend on the Higgs mass $m_H$ through $B_H$. For example, $B_H \sim 1$ for $m_H = 100$ and drops rapidly once the $H \rightarrow W^+W^-$ decay channel opens (note that for $m_H > 2m_t$ the Higgs can decay to a $t\bar{t}$ pair followed by $t\bar{t} \rightarrow b\bar{b} + X$, and $B_H$ is therefore non-negligible in this mass range). In Fig. 1 and in what follows we use the conservative value $m_H = 100$ GeV for which $B_H$ is $\sim 1$. This choice has almost no effect on $\sigma_B^{(1)}$ which is dominated by $\sigma_{t\bar{t}}$, $\sigma_{tt}$ and $\sigma_{ZZ}$. In contrast $\sigma_B^{(3)}$ is sensitive to $m_H$, and in some cases may be up to two times smaller if $m_H > 2m_W$. We also plot the SM CKM-suppressed irreducible backgrounds $e^+e^- \rightarrow W^+W^- \rightarrow 1b-$jet $+ X$ (see Fig. 1(a)) and $e^+e^- \rightarrow t\bar{t} \rightarrow b\bar{b}W^+W^- \rightarrow 3b-$jets $+ X$ (see Fig. 1(b)) for a 500 GeV $e^+e^-$ collider. Clearly, these cross-sections are much smaller than the corresponding reducible $b_P = 1$ cross-sections.

As expected, the background for $n = 1$ $b_P = -1$ processes is typically about one to two orders of magnitudes larger than that for $n = 3$ processes due to the large $e^+e^- \rightarrow b\bar{b}$, $t\bar{t}$ cross-sections which do not contribute to $\sigma_B^{(3)}$. Thus $b$-tagging generally (but not necessarily as will be shown below) produces a more sensitive probe of new physics in $e^+e^-$ colliders for $b_P = -1$ processes with $n = 3$. The figures also show that $\sigma_B^{(1,3)} \rightarrow 0$ as $\epsilon_b \rightarrow 1$ in which case there is no background to $b_P = -1$ processes (neglecting the very small SM irreducible $b_P = -1$ contributions) since in this limit it would be possible to identify a $b$-jet with 100% efficiency.

Let us now consider some specific examples of new physics with $b_P = -1$. To be as general as possible, we take a model independent approach in which
we investigate the limits that can be placed on the scale $\Lambda$ of a new short-distance theory which generates flavor violations by analyzing $b_P$-violating effective operators which give rise to $n = 1$ and $n = 3$ processes. The low-energy effective Lagrangian generated by the underlying theory, which respects the SM $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge invariance, with new flavor dynamics [8] can be written in the form (to lowest order in $1/\Lambda$)

$$\mathcal{L}_{eff} = \frac{1}{\Lambda^2} \sum_i \alpha_i \mathcal{O}_i ,$$

where $\mathcal{O}_i$ are mass dimension 6 effective operators and $\alpha_i$ are coefficients expected to be of $\mathcal{O}(1)$ [8].

Figure 1: Reducible $b_P$-even SM background (in pb) to $b_P$-odd signals with (a) $n = 1$ and (b) $n = 3$ $b$-jets in the final state, as a function of the $b$-tagging efficiency $\epsilon_b$, and for $e^+e^+$ colliders with $\sqrt{s} = 200$, 500 and 1000 GeV in descending order, respectively (solid lines). Also shown are the SM $b_P$-odd irreducible cross-sections (dashed lines) for (a) $n = 1$ and (b) $n = 3$ signals. The Higgs mass was set to 100 GeV.
As a concrete example we first consider $n = 1$ processes generated by the following four-Fermi operator

$$\mathcal{O}^{(1)}_{\ell q} = (\bar{\ell} \gamma^{\mu} \ell) (\bar{q}_i \gamma_{\mu} q_j), \quad (8)$$

where $\ell$ and $q$ are the SM left-handed lepton and quark $SU(2)_L$ doublets and $i, j = 1, 2$ or 3 label the generation. This operator gives rise to a new four-Fermi $e^+e^-t\bar{c}$ and $e^+e^-b\bar{s}$ vertices (and their charged conjugates). It can be generated, for example, by an exchange of a heavy vector-boson in the underlying theory (see [8]). The corresponding cross-sections are [9]:

$$\sigma_{tc} \equiv \sigma(e^+e^- \rightarrow t\bar{c} + \bar{t}c) = \frac{s}{\Lambda^4} \frac{|\alpha^{(1)}_{\ell q}|^2 \beta_t^2 (3 + \beta_t)}{4\pi (1 + \beta_t)^3}, \quad (9)$$

where $\beta_t = (s - m_t^2)/(s + m_t^2)$. $\sigma_{bs} \equiv \sigma(e^+e^- \rightarrow b\bar{s} + \bar{b}s)$ is obtained from $\sigma_{tc}$ by taking $\beta_t = 1$.

Since we are interested here in the inclusive one $b$-jet production $e^+e^- \rightarrow 1b-$jet + $X$ (i.e., $n = 1$) our observable signal is $\sigma^{(1)}_S = P_1^S \times (\sigma_{tc} + \sigma_{bs})$, where we included the $b$-tagging efficiency $P_1^S = \epsilon_b$ for identifying the only $b$-jet in the final state.

In Table 1 we give the $3\sigma$ limits (in case such a signal is not observed) that can be placed on the scale $\Lambda$, assuming that $|\alpha^{(1)}_{\ell q}| = 1$ and based only on our $b$-tagging method. We thus require $\left(\frac{\sigma^{(1)}_S}{\sqrt{\sigma^{(1)}_B}}\right) \times \sqrt{L} \geq 3$, where $L$ is the yearly integrated luminosity of the $e^+e^-$ collider. We consider three $e^+e^-$ collider scenarios: LEP2 with c.m. energy of $\sqrt{s} = 200$ GeV, $L = 2.5$ fb$^{-1}$, a next linear collider (NLC) with $\sqrt{s} = 500$ GeV and $L = 75$ fb$^{-1}$ and a NLC with $\sqrt{s} = 1000$ GeV and $L = 200$ fb$^{-1}$. We have set $m_H = 100$ GeV, but as mentioned before, the limits in Table 1 are insensitive to the choice of the SM Higgs mass.

Table 1: $3\sigma$ limits on $\Lambda$, the scale of the new $b_p = -1$ physics which generates the four-Fermi effective operator in (8). Limits are obtained for
\( \epsilon_b = 50\%, \, 60\% \) and 70\%, by considering the corresponding inclusive 1b-jet signal \( e^+e^- \rightarrow 1b - \text{jet} + X \) without any cuts (see text). The limits are given for \( |\alpha^{(1)}_{\ell q}| = 1, \, m_h = 100 \text{ GeV} \) and for three accelerator scenarios; \( \sqrt{s} = 200, \, 500 \) and 1000 GeV with luminosities 2.5, 75 and 200 \( fb^{-1} \), respectively.

<table>
<thead>
<tr>
<th>( \sqrt{s} )</th>
<th>( L )</th>
<th>( \epsilon_b = 50% )</th>
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</thead>
<tbody>
<tr>
<td>200 GeV</td>
<td>2.5 ( fb^{-1} )</td>
<td>1.4 TeV</td>
<td>1.5 TeV</td>
<td>1.6 TeV</td>
</tr>
<tr>
<td>500 GeV</td>
<td>75 ( fb^{-1} )</td>
<td>5.0 TeV</td>
<td>5.2 TeV</td>
<td>5.5 TeV</td>
</tr>
<tr>
<td>1000 GeV</td>
<td>200 ( fb^{-1} )</td>
<td>9.5 TeV</td>
<td>10.0 TeV</td>
<td>10.7 TeV</td>
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We see from Table 1 that the obtainable limits on the scale of such new physics are typically \( \Lambda > 8 - 10\times \sqrt{s} \), where \( \sqrt{s} \) is the c.m. energy in the given collider. For example, a 3\( \sigma \) limit (on such a \( b_P = -1 \) four-Fermi operator) of \( \Lambda > 1.5 \text{ TeV} \) is obtainable already at LEP2 energies using our \( b \)-jets counting method with 60\% \( b \)-tagging efficiency. These limits are comparable with those obtained using other methods, see [10] and references therein.

Next we consider a \( b_P = -1 \) effective operator that can give rise to \( n = 3 \) signals in a leptonic collider

\[ O_{\phi\phi} = \left( \phi^\dagger \epsilon^D \phi \right) \left( \bar{u}_i \gamma_\mu d_j \right), \quad (10) \]

where \( u_i \) and \( d_j \) are SM right-handed up and down quark singlets \( (i, j = 1, 2 \) or 3 is a generation index) and \( \epsilon_{12} = -\epsilon_{21} = 1 \) [8]. This operator alters the \( Wcb \) vertex with a new right handed coupling as follows

\[ L_{Wcb} = \frac{g}{\sqrt{2}} \gamma_\mu \left( V_{cb} L + \frac{v^2}{\Lambda^2} \alpha_{\phi\phi} R \right), \quad (11) \]

where \( v \) is the SM Higgs vacuum expectation value \( (v = 246 \text{ GeV}) \), \( L(R) = (1 - (+)\gamma_5)/2 \) and, again, \( \alpha_{\phi\phi} \) is naturally of \( O(1) \). In fact, this case is of
particular interest since $\mathcal{L}_{Wcb}$ gives rise to $b_P = -1$ processes with both $n = 1$ and $n = 3$. These are: $e^+ e^- \to W^+ W^- \to b + X$ or $\bar{b} + X$ and $e^+ e^- \to t\bar{t} \to b\bar{b}b + X$ or $b\bar{b}\bar{b} + X$, respectively. The corresponding cross-sections can be calculated in the narrow width approximation in which the entire dependence on the new right-handed coupling comes from the branching ratios of $W^+ \to c\bar{b}$ or $W^- \to \bar{c}b$ (we expand to lowest order in the new right-handed $Wcb$ coupling):

$$
\sigma_{Wcb}^{(1)} = \sigma(e^+ e^- \to W^+ W^- \to 1b\text{-jet} + X) \\
\simeq \sigma(e^+ e^- \to W^+ W^-) \times [Br(W^+ \to c\bar{b}) + Br(W^- \to \bar{c}b)] ,
$$

$$
\sigma_{Wcb}^{(3)} = \sigma(e^+ e^- \to t\bar{t} \to 3b\text{-jets} + X) \\
\simeq \sigma(e^+ e^- \to t\bar{t}) \times [Br(W^+ \to c\bar{b}) + Br(W^- \to \bar{c}b)] ,
$$

where

$$
Br(W^+ \to c\bar{b}) = Br(W^- \to \bar{c}b) = 3 \left( |V_{cb}|^2 + \frac{v^4}{\Lambda^4} |\alpha_{\phi\phi}|^2 \right) \frac{G_F m_W^3}{\sqrt{2} \, 6\pi} ,
$$

and we took a branching ratio = 1 for the second $W$-boson in (12).

In Table 2 we give the $3\sigma$ limits on the scale $\Lambda$ of such new physics which gives rise to a new right-handed $Wcb$ coupling, assuming $|\alpha_{\phi\phi}| = 1$ and based only on the proposed $b$-tagging method. We consider two $e^+ e^-$ collider scenarios; LEP2 with a c.m. energy of 200 GeV and an integrated luminosity of $L = 2.5$ fb$^{-1}$ and a 500 GeV NLC with $L = 75$ fb$^{-1}$; again we set $m_H = 100$ GeV. The limits obtained using the $n = 3$ case are given in parenthesis only for a 500 GeV NLC (a $t\bar{t}$ cannot be produced in LEP2). We find that the limits on $\Lambda$ cannot be improved in a 1000 GeV NLC with $L = 200$ fb$^{-1}$. We see that, typically, the limits are $\Lambda > 800$ GeV in LEP2 and $\Lambda > 1200$ GeV in a 500 GeV $e^+ e^-$ collider.

The above numbers can be compared, for example, with the bound obtained from the measurement of the $W$-boson total width. Using $\Gamma_W =$
2.06 ± 0.06 GeV [1], and assuming that the central value agrees with the SM prediction, we find that the uncertainty in \( \Gamma_W \) places the limit \( \Lambda \sim > 450 \ (1\sigma) \), 340 \ (3\sigma) \) GeV on the scale of the new right-handed \( Wcb \) operator. This is about 2.5 times smaller than the limits given in Table 2 already at LEP2 energies.

Note that the \( \sigma^{(3)}_{Wcb} \) also contributes to the inclusive \( n = 1 \) signal whenever two \( b \)-jets are missed. This effect, though small, is included in the results.

Table 2: \( 3\sigma \) limits on \( \Lambda \) the scale of the \( b_P = -1 \) new physics which generates the effective operator (10). The limits are obtained using \( b \)-tagging efficiency only, with \( \epsilon_b = 50\% \), \( 60\% \) and \( 70\% \), by considering the inclusive \( 1b \)-jet signal \( e^+e^- \rightarrow 1b \)-jet + \( X \) for the case \( e^+e^- \rightarrow W^+W^- \rightarrow 1b \)-jet + \( X \) (see text) and the inclusive \( 3b \)-jet signal \( e^+e^- \rightarrow 3b \)-jet + \( X \) for the case \( e^+e^- \rightarrow t\bar{t} \rightarrow 3b \)-jets + \( X \) (in parenthesis, for a 500 GeV collider). See also Table 1.

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</tr>
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<td>75 fb(^{-1} )</td>
<td>1100 (700) GeV</td>
<td>1200 (800) GeV</td>
<td>1200 (900) GeV</td>
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It is interesting to note that in spite of the much smaller background to the \( n = 3 \) process (see Fig. 1), the sensitivity of a 500 GeV collider to \( \Lambda \) is slightly higher for the \( n = 1 \) processes than for \( n = 3 \). This is not the general case and results from the large \( W^+W^- \) production rate (more than an order of magnitude larger than the \( t\bar{t} \) rate) for this type of collider. We also note in passing that, assuming \( \alpha_{\phi\phi} = 0 \), \( \Lambda \sim 1.2 \) TeV corresponds to \( |V_{cb}| \sim 0.04 \)
which is within the existing experimentally measured value. Thus, $V_{cb}$ effects in high energy $e^+e^-$ colliders may be studied through our $b$-tagging method.

Before we summarize we wish to note that the following issues need further investigation:

- The proposed $b$-jet counting method can be used to constrain specific models containing $b_P = -1$ interactions. For example, supersymmetric models without R-parity (see e.g., [11]), or multi-Higgs doublet models without natural flavor conservation (see e.g., [12]) can give rise to $t\bar{c}, t\bar{u}$ (or $b\bar{s}, b\bar{d}$) production in leptonic colliders, which lead to $n = 1$, $b_P$-odd signals. These type of models can also give rise to $n = 3$, $b_P = -1$ signals, e.g., via top quark decays. Moreover, flavor violation in the scalar sector of an R-parity conserving supersymmetry may also lead to $b_P = -1$ signals in leptonic colliders.

- In leptonic colliders with c.m. energies $\gtrsim 1.5$ TeV, $t$-channel vector-boson fusion processes become important (see e.g., [13]). At such energy scales, the SM $b_P = 1$ reducible background needs to be reevaluated. At the same time, the $V_1V_2$-fusion processes ($V_{1,2} = \gamma, Z$ or $W$) give rise to a variety of new possible $b_P = -1$ signals from new flavor physics interactions (see e.g., S. Bar-Shalom et al. in [12]).

- In hadronic colliders there are additional problems due to the $b$-quark content of the protons (essentially from gluon splittings). These initial-state $b$-quarks may produce significant irreducible SM $b_P$-odd signals, e.g., $b + \text{gluon} \rightarrow b + \text{gluon}$, which need to be taken into account. However, at the Tevatron 2 TeV $p\bar{p}$ collider, due to its relatively low c.m. energy, the $b$-quark content of the protons is small and we expect that the $b_P = 1$ irreducible background considered in this letter (now generated in $q\bar{q}$ annihilation) will still dominate. In this situation new
\[ b_P = -1 \] signals can be also studied at the Tevatron using our \(b\)-tagging method without much difficulty.

To summarize, we have shown that \(b\)-tagging alone can be used to efficiently probe physics beyond the SM in inclusive \(b\)-jet production processes \(\ell^+\ell^- \rightarrow nb + X\) \((n\) is the number of \(b\)-jets in the final state\). We suggested that these reactions can be characterized through the use of an approximate quantum number \(b_P = (-1)^n\) which we called \(b\)-Parity. Due to small off-diagonal CKM matrix elements, \(b_P\) is conserved within the SM to very good accuracy; it follows that the SM contributions to the above reactions are \(b_P\)-even. The only limitation in using this counting procedure for probing new flavor physics is due to a reduced \(b\)-tagging efficiency \(\epsilon_b\). Despite the presence of this (reducible) background, we showed that our method is sensitive enough to provide very useful limits on new flavor physics in a variety of scenarios (of which two examples are provided) using realistic values for \(\epsilon_b\).

We thank David Atwood for suggesting the name \(b\)-Parity. This research was supported in part by US DOE contract number DE-FG03-94ER40837(UCR).
References


