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# Associated production of $\Upsilon$ and weak gauge bosons in hadron colliders

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We calculate the rate of production of $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ in hadron colliders. We find the cross sections for $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ to be roughly 0.45 pb and 0.15 pb at the Fermilab Tevatron and roughly 4 pb and 2 pb at the CERN Large Hadron Collider (LHC). The dominant production mechanism involves the binding of a color-octet $b\bar{b}$ pair into a $P$-wave bottomonium state which subsequently decays into $\Upsilon$. The purely leptonic decay modes of $\Upsilon$, $W^{\pm}$, and $Z^0$ provide signatures with small backgrounds. These events may be observable in run II at the Tevatron, and they should certainly be observable at the LHC. [S0556-2821(99)50317-2]

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In high energy collider experiments, the particles with the cleanest signatures are those with purely leptonic decay modes. They include the charged weak gauge bosons $W^{\pm}$, which decay via $W\rightarrow l\nu$ where $l$ is an electron or muon, and the neutral weak gauge boson $Z^0$, which decays via $Z^0\rightarrow l^+l^-$. They also include the $J^{PC}=1^{--}$ quarkonium states, such as the charmonium state $J/\psi$ and the bottomonium state $\Upsilon$, which decay into $l^+l^-$. Purely leptonic decay modes are particularly useful in hadron colliders, because they provide an enormous suppression of the background. Events involving particles with such decay modes can be employed both as a probe of the production process and as a lamppost under which to search for new physics. In this Rapid Communication, we calculate the rate in the standard model for the associated production of the bottomonium state $\Upsilon$ with $W^{\pm}$ or $Z^0$ at the Fermilab Tevatron proton-antiproton collider and at the CERN Large Hadron Collider (LHC). We find that the dominant production mechanism for both $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ involves the binding of a color-octet $b\bar{b}$ pair into a $P$-wave bottomonium state which subsequently decays into $\Upsilon$. The purely leptonic decays of the $\Upsilon$, $W^{\pm}$, and $Z^0$ provide clean signatures for these events. The cross sections for $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ may be too small for these events to be observed at the Tevatron, but they should certainly be seen at the LHC.

The associated production of $W^{\pm}$ or $Z^0$ and a $J/\psi$ has been studied previously, but only for $J/\psi$'s with large transverse momentum, where the process is dominated by gluon fragmentation into $J/\psi$ [1]. The differential cross sections for $W^{\pm}+J/\psi$ and $Z^0+J/\psi$ peak at small transverse momentum of the $J/\psi$. Unfortunately, if the $J/\psi$ is produced with small transverse momentum, one or more of the leptons from its decay is likely to have insufficient energy to be identified. Because the mass of the $\Upsilon$ is larger than that of the $J/\psi$ by a factor of 3, the cross sections for $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ are much smaller than those for $W^{\pm}+J/\psi$ and $Z^0+J/\psi$. However the larger mass of the $\Upsilon$ makes it possible to observe the leptons from its decay, even if the $\Upsilon$ is produced at small transverse momentum. It is therefore possible to measure the total cross sections for $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$.

The production of $\Upsilon$ requires the creation of a $b\bar{b}$ pair at short distances of order $1/m_b$ or smaller. The $b\bar{b}$ pair subsequently forms a color-singlet bound state, which can either be $\Upsilon$ or a higher bottomonium state that decays into $\Upsilon$. The probability of binding depends on the bottomonium state and on the color and angular momentum quantum numbers of the $b\bar{b}$ pair. In the nonrelativistic QCD factorization approach to inclusive quarkonium production [2], these probabilities are parameterized by nonperturbative matrix elements in NRQCD. The matrix elements scale in a definite way with the typical relative velocity $v$ of the bottom quark in bottomonium. Its value is roughly $v^2\approx 1/10$. The NRQCD scaling rules can be used to estimate the order of magnitude of the binding probability, with every factor of $v^2$ corresponding to suppression by an order of magnitude.

We first consider the direct production of $\Upsilon$. The probability of forming an $\Upsilon$ is largest for the $b\bar{b}_1(^3S_1)$ state, which is a color-singlet configuration with angular-momentum quantum numbers $^3S_1$. It can be parametrized by the NRQCD matrix element $\langle\mathcal{O}_1^{\Upsilon}(^3S_1)\rangle$, which is proportional to the square of the wave function at the origin and can be extracted phenomenologically from the leptonic width of the $\Upsilon$: $\langle\mathcal{O}_1^{\Upsilon}(^3S_1)\rangle=7\ \text{GeV}^3$. According to the velocity-scaling rules of NRQCD, the $b\bar{b}$ configurations with the next largest probabilities for binding into an $\Upsilon$ are the color-octet states $b\bar{b}_8(^3S_1)$, $b\bar{b}_8(^1S_0)$, and $b\bar{b}_8(^3P_J)$ [2]. Their binding probabilities are suppressed by $v^4$ relative to $b\bar{b}_1(^3S_1)$. If the cross section for color-octet $b\bar{b}$ pairs was comparable to that for $b\bar{b}_1(^3S_1)$, then the color-octet contributions to the cross section for direct $\Upsilon$ production would be suppressed by two orders of magnitudes. However in hadron colliders, the cross section for color-octet $b\bar{b}$ pairs is much larger than that for $b\bar{b}_1(^3S_1)$, and this can compensate for the smaller binding probability. The cross section for $b\bar{b}_8(^3S_1)$ is particularly large, because it can be produced by the decay of an off-shell gluon. There are as yet no reliable phenomenological determinations of the matrix elements $\langle\mathcal{O}_8^{\Upsilon}(^3S_1)\rangle$, $\langle\mathcal{O}_8^{\Upsilon}(^1S_0)\rangle$, and $\langle\mathcal{O}_8^{\Upsilon}(^3P_0)\rangle$. In their analysis of $\Upsilon$ production at the

 

Tevatron [3], Cho and Leibovich used the value $\langle \mathcal{O}_8^{\Upsilon}(^3S_1)\rangle \approx 0.006\ \text{GeV}^3$, which they obtained by taking a phenomenological value for $\langle \mathcal{O}_8^{J/\psi}(^3S_1)\rangle$ with large error bars and scaling it by $(m_b v_b)^3/(m_c v_c)^3$. They extracted a value for the linear combination $\langle \mathcal{O}_8^{\Upsilon}(^1S_0)\rangle/5+\langle \mathcal{O}_8^{\Upsilon}(^3P_0)\rangle/m_b^2$ by fitting CDF data for bottomonium production [4]. They obtained a central value 0.008 $\text{GeV}^3$ with a statistical uncertainty that was greater than 100%. Thus the uncertainties in the color-octet matrix elements are probably at least an order of magnitude.

We next consider the indirect production of $\Upsilon$ by the binding of the $b\bar{b}$ pair into a higher bottomonium state that subsequently decays into $\Upsilon$. The higher states include $\Upsilon(2S)$ and $\Upsilon(3S)$, which decay into $\Upsilon$ with branching fractions of about 31% and 13%, respectively. Since these states have the same quantum numbers $J^{PC}=1^{--}$ as $\Upsilon$, their NRQCD matrix elements scale in the same way with $v$ and will therefore be comparable in magnitude to those for $\Upsilon$. For example, the binding probabilities of $b\bar{b}_1(^3S_1)$ into $\Upsilon(2S)$ and $\Upsilon(3S)$ are both about 40% smaller than for $\Upsilon$. The effect of the feeddown from higher bottomonium states on the production of $\Upsilon$ from $b\bar{b}_1(^3S_1)$ can be taken into account in the NRQCD factorization formulas by replacing $\langle \mathcal{O}_1^{\Upsilon}(^3S_1)\rangle$ by

$$\sum_H B(H\to\Upsilon X)\langle \mathcal{O}_1^{H}(^3S_1)\rangle = 8\ \text{GeV}^3, \tag{1}$$

where we have added the contributions from $\Upsilon$ (with $B=1$), $\Upsilon(2S)$, and $\Upsilon(3S)$. Including the feeddown from $\Upsilon(2S)$ and $\Upsilon(3S)$ therefore increases the cross section by about 20%.

Bottomonium states with different $J^{PC}$ quantum numbers can have a more dramatic effect on the cross section, because they receive contributions from parton processes with different selection rules. The only process whose binding probability is of the same order in $v$ as $b\bar{b}_1(^3S_1)$ into $\Upsilon(nS)$ is $b\bar{b}_1(^1S_0)$ into $\eta_b(nS)$. By the heavy-quark spin symmetry of NRQCD, the matrix element $\langle \mathcal{O}_1^{\eta_b(nS)}(^3S_1)\rangle$ differs from that of $\langle \mathcal{O}_1^{\Upsilon(nS)}(^1S_0)\rangle$ only by a spin factor of 1/3. Including the contributions from $\eta_b(2S)$ and $\eta_b(3S)$ and assuming that $B(\eta_b(nS)\to\Upsilon X)$ is comparable to the branching fraction for hadronic transitions from $\Upsilon(nS)$ to $\Upsilon$, we obtain the estimate $\Sigma_H B(H\to\Upsilon X)\langle \mathcal{O}_1^{H}(^1S_0)\rangle \approx 0.3\ \text{GeV}^3$.

The binding probabilities for $P$-wave bottomonium states are suppressed by $v^2$ relative to $\Upsilon$. For $\chi_{bJ}(nP)$, the $b\bar{b}$ configurations with binding probabilities suppressed by only $v^2$ are $b\bar{b}_1(^3P_J)$ and $b\bar{b}_8(^3S_1)$. By heavy-quark spin symmetry, the color-singlet matrix elements $\langle \mathcal{O}_1^{\chi_{bJ}(nP)}(^3P_J)\rangle$ differ for $J=0,1,2$ only by a spin factor of $2J+1$. Their values can be determined from the wave function in nonrelativistic potential models. We adopt the values used in the analysis of Ref. [3]. Weighting the matrix elements for $\chi_{bJ}(1P)$ and $\chi_{bJ}(2P)$ by their measured branching fractions into $\Upsilon$, we obtain $\Sigma_H B(H\to\Upsilon X)\langle \mathcal{O}_1^{H}(^3P_J)\rangle/m_b^2 \approx (0.003,0.2,0.2)\ \text{GeV}^3$ for $J=(0,1,2)$. We have included a factor of $1/m_b^2$ in the $P$-wave matrix elements so that they have the same dimensions as $S$-wave matrix elements. The relative magnitudes of the matrix elements can then be interpreted as relative binding probabilities. The color-octet matrix elements $\langle \mathcal{O}_8^{\chi_{bJ}(nP)}(^3S_1)\rangle$ differ for $J=0,1,2$ only by a spin factor of $2J+1$. In the analysis of $\Upsilon$ production at the Tevatron by Cho and Leibovich [3], these matrix elements were determined by fitting the CDF data on bottomonium production [4]. Weighting the matrix elements by their measured branching fractions into $\Upsilon$, we obtain

$$\sum_H B(H\to\Upsilon X)\langle \mathcal{O}_8^{H}(^3S_1)\rangle \approx 0.4\ \text{GeV}^3. \tag{2}$$

We have included the contributions from $\chi_{bJ}(1P)$ with $J=1,2$ and from $\chi_{bJ}(2P)$ with $J=0,1,2$. The $1P$ states account for about 70% of the total. The value (2) is more than an order of magnitude larger than the estimate by Cho and Leibovich of the matrix element $\langle \mathcal{O}_8^{\Upsilon}(^3S_1)\rangle$ that contributes to direct $\Upsilon$ production. The uncertainty in (2) is probably at least a factor of 3, because the analysis of $\Upsilon$ production has substantial theoretical uncertainties. The weighted matrix element in (2) is smaller than the corresponding color-singlet term in (1) by an order of magnitude. This smaller binding probability is more than compensated by the much larger cross section for $b\bar{b}_8(^3S_1)$ in hadron colliders. Thus the binding of $b\bar{b}_8(^3S_1)$ into $\chi_{bJ}$ followed by its decay into $\Upsilon$ gives a large contribution to the total $\Upsilon$ cross section.

We now consider the production of $W^{\pm}+\Upsilon$ in a hadron collider. This process proceeds through parton collisions that produce final states including $W+b\bar{b}$. The relative importance of the various contributions depends on the magnitude of the NRQCD matrix elements and on the size of the parton cross sections. The only parton cross sections that can be calculated relatively easy are those that can be taken into account via $2\to3$ parton processes of the form $ij\to W+b\bar{b}$. Fortunately, they account for the largest contributions. The $2\to3$ parton processes that produce $W^{+}+b\bar{b}$ are

W1: $u\bar{d},c\bar{s}\to W^{+}+b\bar{b}_8(^3S_1)$ via the decay of a virtual gluon into $b\bar{b}$,

W2: $c\bar{b}\to W^{+}+b\bar{b}_1(^3S_1,{}^1S_0,{}^3P_J)$ involving the exchange of a virtual gluon,

W3: $u\bar{d},c\bar{s}\to W^{+}+b\bar{b}_1(^3S_1)$ via the decay of a virtual photon into $b\bar{b}$.

We have omitted processes that are suppressed by a tiny Kobayashi-Maskawa factor. Process W1 is the only color-octet contribution with a cross section of order $\alpha_s^2\alpha_w$. All other color-octet contributions are suppressed by an additional factor of $\alpha_s$. Color-singlet contributions involving QCD interactions are suppressed by an additional factor of $\alpha_s^2$. One such process, $gg\to W^{+}+b\bar{b}_1(^3S_1)+b\bar{c}$, which has a cross section of order $\alpha_s^4\alpha_w$, is dominated by the splitting of the colliding gluons into a collinear $b\bar{b}$ pair and a collinear $c\bar{c}$ pair. It can be taken into account through process W2. We expect this process to provide a reasonable estimate of the color-singlet contributions involving QCD interactions, because it is the only such process that is enhanced by a factor of $\log(M_W/m_b)\log(M_W/m_c)$. Process W3 is a color-singlet contribution that involves only electroweak interactions and has a cross section of order $\alpha^2\alpha_w$.

We proceed to calculate the cross section for production of $W^\pm+\Upsilon$ at the Tevatron. We use the leading order Glück-Rey-Vogt 1994 (GRV-94) parton distributions, and set the factorization and renormalization scales equal. We neglect the masses of the $u$, $d$, $s$, and $c$ quarks, and we use the value 4.7 GeV for the $b$ quark mass. The cross section in $p\bar{p}$ collisions at center-of-mass energy 1.8 TeV is

$$\sum_{\pm}\sigma(W^\pm+\Upsilon)=1.6\ \text{fb}\ \frac{\sum B\langle\mathcal{O}_1(^3S_1)\rangle}{8\ \text{GeV}^3}+450\ \text{fb}\ \frac{\sum B\langle\mathcal{O}_8(^3S_1)\rangle}{0.4\ \text{GeV}^3}. \tag{3}$$

We have omitted the contributions from the color-singlet $^1S_0$ and $^3P_J$ matrix elements, because they are several orders of magnitude smaller than the color-singlet $^3S_1$ term. The color-singlet $^3S_1$ term is completely dominated by the purely electroweak process W3. However it is overwhelmed by the color-octet $^3S_1$ contribution from process W1. While the color-singlet matrix element is accurately normalized by the leptonic decay width of the Υ, there is a large uncertainty in the normalization of the color-octet matrix element. The theoretical error due to higher order perturbative corrections can be estimated by varying the factorization and renormalization scales by a factor of 2 and is roughly 35%. The theoretical error due to the uncertainty in the $b$ quark mass can be estimated by varying $m_b$ by 0.25 GeV and is roughly 25%.

To obtain the cross section in the purely leptonic decay channels, we multiply by the branching fractions 5% for $\Upsilon\to l^+l^-$ and 22% for $W\to l\nu$. Given the integrated luminosity of 110 $\text{pb}^{-1}$ in run I of the Tevatron, there should be about 0.5 events in the leptonic decay channels. To estimate the number of events that would actually be detected, we must multiply by the detector acceptances and efficiencies. The product of the acceptances and efficiencies for the Collider Detector at Fermilab (CDF) was about 26% for $W\to e\nu$ [5] and about 12% for $\Upsilon\to\mu^+\mu^-$ [4]. Thus after allowing for acceptances and efficiencies, the cross section is about two orders of magnitude too small to be observable in the data from run I.

We next consider the production of $Z^0+\Upsilon$ in a hadron collider. The $2\to3$ parton processes that produce $Z^0+b\bar{b}$ are

Z1: $u\bar{u},d\bar{d},s\bar{s},c\bar{c}\to Z^0+b\bar{b}_8(^3S_1,{}^1S_0,{}^3P_J)$,
Z2: $gg\to Z^0+b\bar{b}_8(^3S_1,{}^1S_0,{}^3P_J)$,
Z3: $gg\to Z^0+b\bar{b}_1(^3S_1,{}^1S_0,{}^3P_J)$,
Z4: $b\bar{b}\to Z^0+b\bar{b}_1(^3S_1,{}^1S_0,{}^3P_J)$,
Z5: $u\bar{u},d\bar{d},s\bar{s},c\bar{c}\to Z^0+b\bar{b}_1(^3S_1)$ via the decay of a virtual photon into $b\bar{b}$.

Processes Z1 and Z2 are color-octet contributions with cross sections of order $\alpha_s^2\alpha_w$. Process Z3 is the only color-singlet contribution with a cross section of order $\alpha_s^2\alpha_w$. Another color-singlet contribution $gg\to Z^0+b\bar{b}_1(^3S_1)+b\bar{b}$, which has a cross section of order $\alpha_s^4\alpha_w$, can be taken into account through process Z4. Finally, process Z5 is an electroweak color-singlet contribution with a cross section of order $\alpha^2\alpha_w$.

The cross section for production of $Z^0+\Upsilon$ in $p\bar{p}$ collisions at the Tevatron with center-of-mass energy 1.8 TeV is

$$\sigma(Z^0+\Upsilon)=8.5\ \text{fb}\ \frac{\sum B\langle\mathcal{O}_1(^3S_1)\rangle}{8\ \text{GeV}^3}+150\ \text{fb}\ \frac{\sum B\langle\mathcal{O}_8(^3S_1)\rangle}{0.4\ \text{GeV}^3}. \tag{4}$$

We have calculated all contributions to direct Υ production with binding probabilities of order $v^4$ or larger and all indirect contributions involving binding probabilities of order $v^2$ or larger. All those terms that are not shown explicitly in (4) give contributions that are more than an order of magnitude smaller than the color-singlet $^3S_1$ term. The color-singlet $^3S_1$ contribution is dominated by processes Z3 and Z5, which contribute about 85% and 15%, respectively. However it is overwhelmed by the color-octet $^3S_1$ term, which is dominated by process Z1. While the color-singlet matrix element is accurately normalized, there is a large uncertainty in the normalization of the color-octet matrix element. We estimate the theoretical errors from higher order perturbative corrections and from the uncertainty in the $b$ quark mass to be roughly 35% and 25%, respectively.

To obtain the cross section in the purely leptonic decay channels, we must multiply by the branching fractions 5% for $\Upsilon\to l^+l^-$ and 7% for $Z^0\to l^+l^-$. Given the integrated luminosity of 110 $\text{pb}^{-1}$ in run I of the Tevatron, there should be about 0.06 events in the leptonic decay channels. The product of the acceptances and efficiencies of the CDF detector was about 30% for $Z^0\to e^+e^-$ [5] and about 12% for $\Upsilon\to\mu^+\mu^-$ [4]. Thus after allowing for acceptances and efficiencies, the cross section is several orders of magnitude too small to be observable in the data from run I.

In run II of the Tevatron, the increase of the center-of-mass energy to 2.0 TeV will increase the cross sections for $W^\pm+\Upsilon$ and $Z^0+\Upsilon$ by about 10%. With an integrated luminosity of 2000 $\text{pb}^{-1}$, the number of events in the purely leptonic decay channels should be about 10 for $W^\pm+\Upsilon$ and about 2 for $Z^0+\Upsilon$. The upgrades of the CDF and DØ detectors should increase the acceptances and efficiencies for observing these events. Given that the uncertainties in the cross sections are at least a factor of 3, there is a possibility that these events could be observed in run II. An observation of these events at a much larger rate than predicted could be evidence for a heavy particle that has a substantial branching fraction into $W^\pm+b\bar{b}$ or $Z^0+b\bar{b}$.

The cross sections for production of $W^\pm+\Upsilon$ and $Z^0+\Upsilon$ in $pp$ collisions at the LHC with center-of-mass energy 14 TeV is

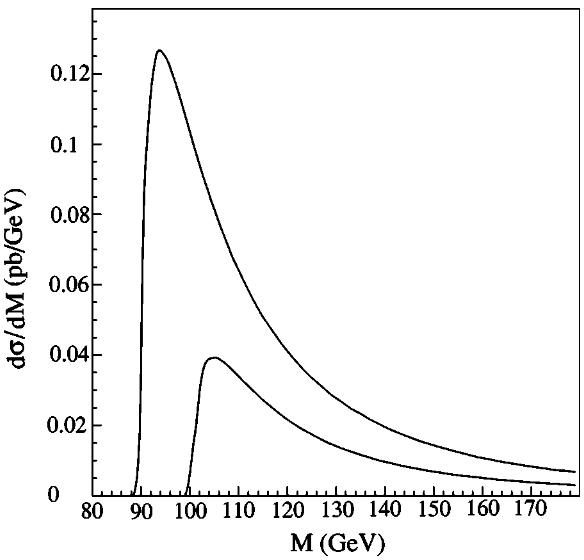


FIG. 1. The invariant mass distributions $d\sigma/dM_{W\Upsilon}$ (upper curve) and $d\sigma/dM_{Z\Upsilon}$ (lower curve) for the production of $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ in $pp$ collisions at center-of-mass energy 14 TeV.

$$\sum_{\pm} \sigma(W^{\pm}+\Upsilon)=10 \;\text{fb}\; \frac{\sum B\langle\mathcal{O}_1(^3S_1)\rangle}{8 \;\text{GeV}^3} + 4000 \;\text{fb}\; \frac{\sum B\langle\mathcal{O}_8(^3S_1)\rangle}{0.4 \;\text{GeV}^3}, \tag{5}$$

$$\sigma(Z^0+\Upsilon)=500 \;\text{fb}\; \frac{\sum B\langle\mathcal{O}_1(^3S_1)\rangle}{8 \;\text{GeV}^3} + 1300 \;\text{fb}\; \frac{\sum B\langle\mathcal{O}_8(^3S_1)\rangle}{0.4 \;\text{GeV}^3}. \tag{6}$$

With an integrated luminosity of 10 $\text{fb}^{-1}$, the number of events in the purely leptonic decay channels should be about 440 for $W^{\pm}+\Upsilon$ and about 70 for $Z^0+\Upsilon$. Even after allowing for detector acceptances and efficiencies, there should be enough events to make these processes observable. In Figure 1, we plot the invariant mass distributions $d\sigma/dM_{W\Upsilon}$ (summed over $W^{\pm}$) and $d\sigma/dM_{Z\Upsilon}$. They peak at only a few GeV above the thresholds of 89.8 GeV for $W^{\pm}+\Upsilon$ and 100.6 GeV for $Z^0+\Upsilon$.

The cross sections for $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ are both dominated by the $\langle\mathcal{O}_8(^3S_1)\rangle$ term. The uncertainty in the value of $\Sigma B\langle\mathcal{O}_8(^3S_1)\rangle$ leads to a substantial uncertainty in the normalizations of the cross sections. However this uncertainty cancels in the ratio $\Sigma_{\pm}\sigma(W^{\pm}+\Upsilon)/\sigma(Z^0+\Upsilon)$, which is predicted to be about 3 at the Tevatron and about 2 at the LHC. If a significant deviation from this prediction is observed at the LHC, it might be evidence for an additional contribution from a heavy particle that decays into $W^{\pm}+b\bar{b}$ or $Z^0+b\bar{b}$.

The production of $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ provide lampposts under which one can look for new physics. The most promising possibility is to search for a charged Higgs boson via the decay $H^+\rightarrow W^+ + \Upsilon$. The decay rate of the charged Higgs boson into $W+b\bar{b}$ is enhanced by the Yukawa coupling of the Higgs boson to a virtual top quark. If the mass of the charged Higgs boson is in the range between 140 GeV and $t\bar{b}$ threshold and if the Higgs mixing parameter $\tan\beta$ is small, then $W+b\bar{b}$ may be the largest single decay mode [6]. The decay rate of the charged Higgs boson into $W+\Upsilon$ was first calculated by Grifols, Gunion, and Mendez [7]. For small $\tan\beta$, the branching fraction $B(H^+\rightarrow W^+ + \Upsilon)$ ranges from about $10^{-4}$ if the Higgs boson mass is just above the $W^+ + \Upsilon$ threshold to about $10^{-3}$ if the Higgs boson mass is just below the $t\bar{b}$ threshold.

In a hadron collider, most of the standard production mechanisms for a charged Higgs boson in the mass range below $t\bar{b}$ threshold involve the production of an additional very massive particle [8]. The standard production mechanisms for $H^+$ include $t\bar{t}$ production followed by the decay $t\rightarrow H^+ b$, $\bar{t}H^+$ production, $W^-H^+$ production, and $H^-H^+$ production. Because of the additional very massive particle, events in which a charged Higgs boson decays into $W+\Upsilon$ will be easily distinguished from $W+\Upsilon$ events produced by standard model processes. There is one potentially significant production process for a charged Higgs boson that could result in events that resemble standard model $W+\Upsilon$ events. That process is $qb\rightarrow q'bH^+$, which proceeds through a Feynman diagram that involves a virtual $W$ and a virtual top quark [9].

We conclude that it should be possible to observe the production of $W^{\pm}+\Upsilon$ and $Z^0+\Upsilon$ from standard model proceses at the LHC. The clean experimental signatures for these events also makes them valuable lampposts under which to search for heavy particles that decay into $W^{\pm}+b\bar{b}$ or $Z^0+b\bar{b}$.

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[1] V. Barger, S. Fleming, and R.J.N. Phillips, Phys. Lett. B **371**, 111 (1996).

[2] G.T. Bodwin, E. Braaten, and G.P. Lepage, Phys. Rev. D **51**, 1125 (1995).

[3] P. Cho and A. Leibovich, Phys. Rev. D **53**, 150 (1996); **53**, 6203 (1996).

[4] CDF Collaboration, F. Abe *et al.*, Phys. Rev. Lett. **75**, 4358 (1995).

[5] CDF Collaboration, F. Abe *et al.*, Phys. Rev. Lett. **76**, 3070 (1996).

[6] S. Moretti and W.J. Stirling, Phys. Lett. B **347**, 291 (1995); **366**, 451 (1996); E. Barradas, J.L. Diaz-Cruz, A. Gutierrez, and A. Rosado, Phys. Rev. D **53**, 1678 (1996); A. Djouadi, J. Kalinowski, and P.M. Zerwas, Z. Phys. C **70**, 435 (1996); E. Ma, D.P. Roy, and J. Wudka, Phys. Rev. Lett. **80**, 1162 (1998).

[7] J.A. Grifols, J.F. Gunion, and A. Mendez, Phys. Lett. B **197**, 266 (1987); R.W. Robinett and L. Weinkauf, Mod. Phys. Lett. A **6**, 1575 (1991).

[8] J.F. Gunion, H.E. Haber, G.L. Kane, and S. Dawson, *The Higgs Hunter's Guide* (Addison-Wesley, Reading, MA, 1990).

[9] S. Moretti and K. Odagiri, Phys. Rev. D **55**, 5627 (1997).