Observation potential for η_b at the Fermilab Tevatron

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We calculate the cross section for η_b production at the Tevatron at next-to-leading-order in the strong coupling and find that more than two millions of η_b's are expected per inverse picobarn of integrated luminosity. We discuss the decay modes into charmed states and suggest that the decays into $D^*D^{(*)}$ mesons might be the most promising channels to observe the η_b in Run II.

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Considering the rich phenomenology of the Υ states, it is quite surprising that spin singlet $b\bar{b}$ states, including the ${}^1S_0$ ground state, have not been observed yet.

Fine and hyperfine splittings of the quarkonia spectra were calculated using phenomenological models for the heavy-quark potential [1,2]. Recent progress both in lattice and in perturbative QCD has allowed the achievement of comparable precisions [3,4]. Various approaches have been successfully adopted to describe the charmonium system and are believed to be even more reliable for the heavier $b\bar{b}$ states, where relativistic effects are less important. Recent determinations lead to a mass splitting between the $\Upsilon(1S)$ and the $\eta_b(1S)$ in the 40-60–MeV range [3–7].

Searches of the η_b have been pursued in various experiments. In e^-e^+ collisions, cross sections for producing spin singlet states are generally small and the signal is rate-limited. This is compensated by a clean environment, which allows for searches in the inclusive decay modes. Following an original suggestion by Godfrey and Rosner [8], the CLEO collaboration has looked for the η_b in the hindered M1 decay of the $\Upsilon(3S)$ and found no signal [9]. At CERN LEP II, the ALEPH collaboration analyzed the $\gamma\gamma$ interactions data, basing their analysis on the QCD prediction of $\Gamma_{\gamma\gamma}$ partial width [10]: no evidence was found in the four- and six-charged-particle decay modes [11].

The situation is exactly reversed in hadron collisions, where the production rates can be large, with millions of events produced at the Tevatron per inverse picobarn of integrated luminosity, yet the intense hadronic activity makes any inclusive analysis unfeasible. In this case it is necessary to identify decay modes that have triggerable signatures and allow for full invariant mass reconstruction of the decaying state. Such exclusive modes are believed to have very small branching ratios. Braaten *et al.* [12] have suggested that the η_b at the Tevatron could be observed via its decay into $J/\psi J/\psi$, with the subsequent leptonic decay of the J/ψ's; experimental efforts have started in this direction [13]. This signature exploits the upgraded abilities of the CDF detector of triggering on soft muons.

The purpose of this note is twofold. First, we provide the prediction for the inclusive cross section of η_b production at the Tevatron, based on the non-relativistic-quantum-chromo-dynamics (NRQCD) calculation at the next-to-leading accuracy in the strong coupling [14]. We find that the cross section is about 4 times larger than the previous available estimate [12]. Second, stemming from an estimate of the corresponding branching ratio, we suggest a new analysis based on the decay of the $\eta_b \to D^*D^{(*)}$ meson pairs. Thanks to secondary vertex trigger capabilities of CDF and to the high resolution achievable in the invariant mass reconstruction (~ 20 MeV), this could turn out to be the most promising search channel at the Tevatron.

According to the NRQCD factorization approach [15], the inclusive cross section for the η_b in $p\bar{p}$ collisions can be written as

$$\sigma(p\bar{p} \to \eta_b + X) = \sum_{i,j} \int dx_1 dx_2 f_{i/p} f_{j/\bar{p}} \hat{\sigma}(ij \to \eta_b), \tag{1}$$

where

$$\hat{\sigma}(ij \to \eta_b) = \sum_n C_n^{ij} \langle 0 | \mathcal{O}_n^{\eta_b} | 0 \rangle. \tag{2}$$

The short-distance coefficients C_n^{ij}, calculable in perturbative QCD, describe the production of a quark-antiquark pair state with quantum number n, while the $\langle 0 | \mathcal{O}_n^{\eta_b} | 0 \rangle$ are the nonperturbative matrix elements that describe the subsequent hadronization of the $b\bar{b}$ pair into the physical η_b state. These matrix elements can be expanded in powers of $v^2 \simeq 0.1$, the relative velocity of heavy quarks in the bound state, so that, to a given accuracy, only a few terms need to be included in the sum over n.

The case of η_b production is particularly simple. Compared to the leading contribution, the ${}^1S_0^{[1]}(b\bar{b})$, the color octect terms, $({}^1S_0^{[8]}, {}^3S_1^{[8]}, {}^1P_1^{[8]})$, are all suppressed by v^4, while the corresponding short-distance coeffi-

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cients start at least at α_S^2 as the singlet contribution [16]. Hence, given that the nonperturbative matrix elements for the singlet production are extracted from Υ decays, there are no unknown parameters (up to corrections of $\mathcal{O}(v^4)$) entering the estimate for η_b production.

The leading-order cross section for $gg \to \eta_b$, Fig. 1(a), is

$$\hat{\sigma}(gg \to \eta_b) = \frac{\pi^3 \alpha_S^2}{36 m_b^3 \hat{s}} \delta\left(1 - \frac{4m_b^2}{\hat{s}}\right) \langle 0 | \mathcal{O}_1^{\eta_b}({}^1S_0) | 0 \rangle. \quad (3)$$

Next-to-leading-order corrections in the strong coupling have been calculated in Refs. [14,17]. These include virtual corrections to the $2 \to 1$ process, Fig. 1(b), and real $2 \to 2$ processes, such as $gg \to \eta_b g$, $gq \to \eta_b q$, Fig. 1(c). The result for the total cross section in $p\bar{p}$ collisions at 1.96 TeV of center-of-mass energy is

$$\sigma(p\bar{p} \to \eta_b + X) = 2.5 \pm 0.3 \mu\mathrm{b}, \quad (4)$$

where we have adopted CTEQ5M1 parton densities, the corresponding two-loop evolution for $\alpha_S(\mu_R)$ and $m_b = 4.75$ GeV. The quoted uncertainty has been estimated by summing (in quadrature) the errors coming from various sources. The first is associated to the choice of the renormalization and factorization scales. Varying them independently in the range $m_b < \mu_R, \mu_F < 4m_b$ gives an uncertainty of 10%. For the nonperturbative matrix element we have used the determination from Υ leptonic decay, $\langle \Upsilon | \mathcal{O}_1({}^3S_1) | \Upsilon \rangle = 3.5 \pm 0.3\ \mathrm{GeV}^3$, which can be related to the η_b production matrix element by using spin-symmetry and vacuum saturation approximation, $\langle 0 | \mathcal{O}_1^{\eta_b}({}^1S_0) | 0 \rangle = \langle \Upsilon | \mathcal{O}_1({}^3S_1) | \Upsilon \rangle$, up to $\mathcal{O}(v^4)$ corrections. Other determinations coming from potential models [18] and lattice calculations [19] fall in the range of the quoted error. Finally, we also included the effect of an uncertainty on the bottom (pole) mass of ± 50 MeV on the cross section. The strong correlation to the nonperturbative matrix element extraction has been exploited to reduce the uncertainty from this source. The result, quoted in Eq. (4), accounts for the "direct" contribution and does not include feed-downs from higher-mass states, such as the h_b.

For the experimental analysis it is important to know the distribution of η_b at small p^T values. This cannot be described accurately by a fixed-order calculation. In this region of the phase space it is necessary to resume higher-order corrections involving soft-gluon radiation. To this aim, we use PYTHIA [20], matched with the exact matrix elements for $ij \to \eta_b k$ describing the high-p^T tail, as outlined in Ref. [21]. This procedure, which has been shown to work well for the analogous process $gg \to H$, has the additional virtue that it can be directly used for simulation in experiments. The results are shown in Fig. 2, where the differential distribution in p^T for the η_b is shown (upper curve). Hadronization and initial k_T effects are not included. The normalization of the inclusive η_b distributions is obtained from the next-to-leading-order (NLO) result, Eq. (4).

As already stated above, the cross section for the η_b in the central region is about 4 times larger than was estimated in Ref. [12] by rescaling the Υ cross section at high transverse momentum. This procedure is expected to provide an underestimate, since it assumes that the p^T spectra for Υ and η_b have a similar shape at small p^T. In fact Υ color-singlet production proceeds at LO through $gg \to \Upsilon g$ and therefore vanishes at $p^T \sim 0$, where the largest part of the η_b's are produced.

With such a large number of events expected, it is interesting to consider in detail the rare exclusive decays that might give triggerable signatures. It is easy to show that direct decays into photons or lepton pairs give either too small branching ratios or very difficult experimental signatures (such as $\eta_b \to \gamma\gamma$ whose branching ratio is $\mathcal{O}(10^{-5})$). In Ref. [12], the branching ratio $\mathrm{Br}(\eta_b \to J/\psi J/\psi)$ was estimated through scaling from the analo-

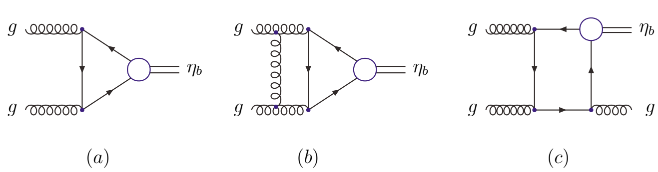

FIG. 1 (color online). Representative Feynman diagrams for η_b hadro-production at LO (a), and virtual (b) and real (c) contributions at NLO.

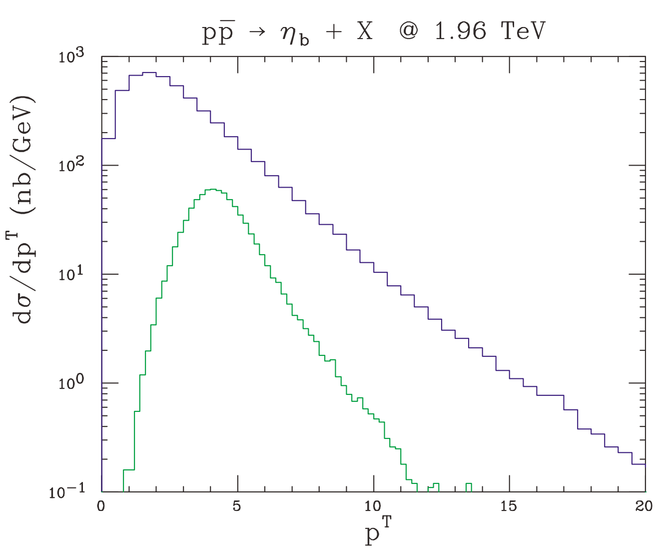

FIG. 2 (color online). Differential cross sections for η_b production at the Tevatron ($p\bar{p}$ at 1.96 TeV). The upper curve is normalized to the NLO total rate and describes the p^T distribution for the η_b. The lower curve is the corresponding distribution for the D mesons coming from the η_b decays, after requiring that they both have $|\eta(D)| < 1$ (no branching ratio included). After the acceptance cut, the rate drops to 15% of the total cross section.

gous $\eta_c \to \phi\phi$, finding a value compatible with $7 \times 10^{-4\pm1}$. This is probably an overestimate. As a simple upper bound, let us consider the inclusive decay rate of the η_b into four-charm states

$$\Gamma(\eta_b \to J/\psi J/\psi) < \Gamma(\eta_b \to c\bar{c}c\bar{c}). \tag{5}$$

The inclusive rate can be calculated at leading-order by considering the four Feynman diagrams, such as the one shown in Fig. 3(b). The result is

$$\mathrm{Br}(\eta_b \to c\bar{c}c\bar{c}) = 1.8^{+2.3}_{-0.8} \times 10^{-5}, \tag{6}$$

where $m_c = 1.45$ GeV, $m_b = 4.75$ GeV, $\alpha_S(2m_b) = 0.182$ and the NLO expression for the inclusive total width [14,22] have been used. The amplitudes have been calculated analytically while the integration over the phase space has been performed numerically. The uncertainty has been estimated by varying the renormalization scale between $m_b < \mu_R < 4m_b$ and the masses in the ± 50 MeV range. The four-charm branching ratio is very sensitive to the value of the charm mass, which dominates its uncertainty. The above result shows that the inclusive rate is already smaller than the lower bound obtained in Ref. [12]; further suppression is expected mainly because many other decay modes to charmed mesons (or other charmonium states) should contribute to the saturation of the inclusive rate. Long distance contributions, which are expected to scale as $\Lambda_{\mathrm{QCD}}/m_b$, cannot be taken quantitatively into account. However, we believe that our result for the inclusive rate is a clear indication that the lower bound of Ref. [12] is likely to be an overestimate.

Given the result in Eq. (6), a comment on the reliability of the estimate based on the scaling from $\mathrm{Br}(\eta_c \to \phi\phi)$ is in order. To this aim we recall that the decay of a scalar $Q\bar{Q}$ meson into two vector states is suppressed in perturbative QCD [23]. A nontrivial check of this selection rule is that the branching ratio $\mathrm{Br}(\eta_b \to J/\psi J/\psi)$ is exactly zero when calculated at LO in the NRQCD double expansion in α_S and v^2. For the same reason one would expect the rate of $\eta_c \to \phi\phi$ to be suppressed, in contradiction with the measured value of about 1%. This entails that some other (nonperturbative) mechanism is responsible for this decay process [24]. Rescaling by $(m_c/m_b)^4$ the

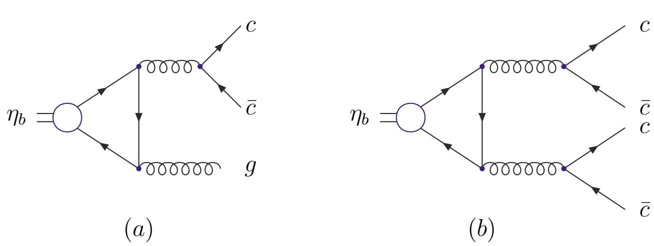

FIG. 3 (color online). Representative Feynman diagrams for the η_b inclusive decay into (a) two-charm (two diagrams) and (b) four-charm (four diagrams) states.

branching ratio of $\eta_c \to \phi\phi$ to obtain the branching ratio of $\eta_b \to J/\psi J/\psi$ amounts to rescaling by the same factor also the effect of nonperturbative or higher-order contributions that are likely to be crucial in determining the η_c decay, but less and less important as we pass from the η_c to the η_b system.

From the phenomenological point of view, Eq. (6) implies $\mathrm{Br}(\eta_b \to \mu^+\mu^-\mu^+\mu^-) < 6 \times 10^{-8}$, which makes a search in this channel very hard. As an alternative, we propose to consider the decays into charmed mesons, $\eta_b \to D^* D^{(*)}$ [28].

A computation of the exclusive decay rates is missing and does not seem to be feasible within the framework of quark models or QCD sum rules. Probably, there is room to face such a problem using lattice techniques. Here we assume that the exclusive decays into $D^* D^{(*)}$ dominate the inclusive rate into charm:

$$\Gamma(\eta_b \to D^* D^{(*)}) \lesssim \Gamma(\eta_b \to c\bar{c} + X), \tag{7}$$

which is a reasonable starting point in absence of a determination of the exclusive channel. We find that the largest contribution to the decay of the η_b into charmed states is given, Fig. 3(a), by

$$\mathrm{Br}(\eta_b \to c\bar{c}g) = 1.5^{+0.8}_{-0.4}\%, \tag{8}$$

where we used the same input parameters and estimated the uncertainties as in the decay into the four-charm states. It can now be argued that not much suppression is expected in passing from the inclusive to the exclusive decays. For example, decay rates into charmonium states should be of the same size $\Upsilon \to J/\psi + X$, *i.e.*, $\mathcal{O}(10^{-3})$ [25]. Emission of extra pions should also be considered, but some of these contributions would be automatically included in the experimental analysis, *e.g.*, $\eta_b \to DD^* \to DD\pi$, as nonresonant diagrams. However, to be conservative, we consider the range $10^{-3} < \mathrm{Br}(\eta_b \to D^* D^{(*)}) < 10^{-2}$ in the following phenomenological analysis.

There are other decay modes leading to a two-charm final state that have not been included in our estimate of the branching ratio. One is the decay of $\eta_b \to g^* g^* \to c\bar{c}$ through a (box) loop, which proceeds at order α_S^4 and is further suppressed by loop factors. Another is the decay $\eta_b \to g^* \to c\bar{c}$, via a $3S_1^{[8]}$ state, which is only, α_S^2 but it is suppressed by the color-octet matrix element $\langle \eta_b | \mathcal{O}_8(^3S_1) | \eta_b \rangle$. Assuming a scaling $\mathcal{O}(v^4)$ for the nonperturbative matrix element, this process gives a non-negligible contribution to the branching ratio, about 5×10^{-3}. However, this result is affected by a large uncertainty and could be much smaller as suggested by studies [26,27] about the size of color octect matrix elements in Υ decays.

Results for the p^T distribution of the D mesons from the η_b decay are shown in Fig. 2 (lower curve). No branching ratio is included, the difference in rate coming only from the requirement that both D mesons be central, $|\eta(D)| < 1$. As expected, the p^T distribution peaks just below $M_{\eta_b}/2$. The efficiency for the geometrical acceptance of the detector is found to be about 15%. By adding the requirement that at least one D meson has $p^T >$ 5 GeV, one is left with only 4% of the total number of events produced. The above efficiency can be folded with our estimate of the branching ratio leading to $10^4 - 10^5$ $D^*D^{(*)}$ triggerable events expected from the η_b decay in 100 pb^{-1} of integrated luminosity at the Tevatron. The final number of reconstructed events will depend on the decay modes of the D and D^* mesons and on the associated experimental efficiencies. We leave this to more detailed experimental studies and only add a few comments. First we note that, according to the arguments outlined above, perturbative QCD predicts the $\eta_b \to D^*D^*$ decay to have a smaller rate with respect to $\eta_b \to D^*D$. In this case, it is reasonable to expect that different charge assignments, such as $D^{*0}D^0$, $D^{*+}D^-$, $D^{*-}D^+$, will occur with the same probability of $\frac{1}{3}$. Finally, we recall that the cleanest signatures have small branching ratios, $\text{Br}(D^0 \to K^-\pi^+) = 3.90\%$ and $\text{Br}(D^+ \to \overline{K}^0\pi^+) = 2.77\%$ [25], leading to a factor of about 10^{-3} drop in the rate if both D mesons are required to decay through these channels. We foresee that sizeable improvements could be achieved by requiring that just one of two D mesons decays through a very clean signature, providing an efficient trigger.

To summarize, we have presented the NLO QCD, prediction for η_b production at the Tevatron, including a resumed result for the p^T spectrum obtained with a dedicated implementation in PYTHIA. The production rate is large, of the order of a few microbarns, and allows for the search of the η_b through rare exclusive decays. We argued that the branching ratio into $J/\psi J/\psi$ is probably too small and we suggested to look for the η_b through its decay into $D^*D^{(*)}$ mesons. Our results indicate that the η_b could be eventually observed during Run II at the Tevatron.

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[1] E. Eichten, K. Gottfried, T. Kinoshita, K. D. Lane, and T.-M. Yan, Phys. Rev. D **21**, 203 (1980).
[2] E. J. Eichten and C. Quigg, Phys. Rev. D **49**, 5845 (1994).
[3] X. Liao and T. Manke, Phys. Rev. D **65**, 074508 (2002).
[4] S. Recksiegel and Y. Sumino, Phys. Lett. B **578**, 369 (2004).
[5] D. Ebert, R. N. Faustov, and V. O. Galkin, Phys. Rev. D **67**, 014027 (2003).
[6] V. Lengyel, Y. Fekete, I. Haysak, and A. Shpenik, Eur. Phys. J. C **21**, 355 (2001).
[7] B. A. Kniehl, A. A. Penin, A. Pineda, V. A. Smirnov, and M. Steinhauser Phys. Rev. Lett. **92**, 242001 (2004).
[8] S. Godfrey and J. L. Rosner, Phys. Rev. D **64**, 074011 (2001).
[9] CLEO Collaboration, A. H. Mahmood *et al.*, hep-ex/0207057.
[10] A. Czarnecki and K. Melnikov, Phys. Lett. B **519**, 212 (2001).
[11] ALEPH Collaboration, A. Heister *et al.*, Phys. Lett. B **530**, 56 (2002).
[12] E. Braaten, S. Fleming, and A. K. Leibovich, Phys. Rev. D **63**, 094006 (2001).
[13] CDF Collaboration, J. Tseng, in *Proceedings of the 5th International Conference on Quark Confinement and the Hadron Spectrum V*, edited by N. Brambilla and G. M. Prosperi (World Scientific, Singapore, 2003).
[14] A. Petrelli, M. Cacciari, M. Greco, F. Maltoni, and M. L. Mangano, Nucl. Phys. B **514**, 245 (1998).
[15] G. T. Bodwin, E. Braaten, and G. P. Lepage, Phys. Rev. D **51**, 1125 (1995).
[16] This is at variance with the J/ψ or Υ production where the C^{gg} coefficient is $\mathcal{O}(\alpha_S^3)$ and therefore color-octet contributions can be important.
[17] J. H. Kuhn and E. Mirkes, Phys. Rev. D **48**, 179 (1993).
[18] E. J. Eichten and C. Quigg, Phys. Rev. D **52**, 1726 (1995).
[19] G. T. Bodwin, D. K. Sinclair, and S. Kim, Phys. Rev. D **65**, 054504 (2002).
[20] T. Sjostrand *et al.*, Comput. Phys. Commun. **135**, 238 (2001).
[21] G. Miu and T. Sjostrand, Phys. Lett. B **449**, 313 (1999).
[22] K. Hagiwara, C. B. Kim, and T. Yoshino, Nucl. Phys. B **177**, 461 (1981).
[23] S. J. Brodsky and G. P. Lepage, Phys. Rev. D **24**, 2848 (1981).
[24] M. Anselmino, F. Murgia, and F. Caruso, Phys. Rev. D **42**, 3218 (1990).
[25] Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002).
[26] F. Maltoni and A. Petrelli, Phys. Rev. D **59**, 074006 (1999).
[27] S. Fleming and A. K. Leibovich, Phys. Rev. Lett. **90**, 032001 (2003).
[28] The decay $\eta_b \to DD$ is forbidden by parity conservation.