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Hadroproduction of $\Upsilon(nS)$ above $B\bar{B}$ Thresholds and Implications for $Y_b(10890)$

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Based on the non-relativistic QCD factorization scheme, we study the hadroproduction of the bottomonium states $\Upsilon(5S)$ and $\Upsilon(6S)$. We argue to search for them in the final states $\Upsilon(1S, 2S, 3S)\pi^+\pi^-$, which are found to have anomalously large production rates at $\Upsilon(5S)$. The enhanced rates for the dipionic transitions in the $\Upsilon(5S)$-energy region could, besides $\Upsilon(5S)$, be ascribed to $Y_b(10890)$, a state reported by the Belle collaboration, which may be interpreted as a tetraquark. The LHC/Tevatron measurements are capable of making a case in favor of or against the existence of $Y_b(10890)$, as demonstrated here. Dalitz analysis of the $\Upsilon(1S, 2S, 3S)\pi^+\pi^-$ states from the $\Upsilon(5S)/Y_b(10890)$ decays also impacts directly on the interpretation of the charged bottomonium-like states, $Z_b(10600)$ and $Z_b(10650)$, discovered by Belle in these puzzling decays.

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As a multi-scale system, heavy-quarkonium states provide a unique laboratory to explore the interplay between perturbative and nonperturbative effects of QCD. Due to the non-relativistic nature, these states allow the application of theoretical tools that can simplify and constrain the analyses of nonperturbative effects. The commonly-accepted method is the non-relativistic QCD (NRQCD) [1] which adopts a factorization ansatz to separate the short-distance and long-distance effects. Since the bottom quark is approximately three times heavier than the charm quark, it is expected that the expansion in $\alpha_s(\mu)$, where μ is a scale of $O(m_b)$, and v^2, with v as the velocity of the heavy quark in the hadron, which also is an NRQCD expansion parameter, converges much faster for the bottomonium states. Consequently, great progress has been made in the past years on the hadronic production of $\Upsilon(1S, 2S, 3S)$ [2]. On the experimental side the production rates and polarization have been measured at the Tevatron [3–5] and at the LHC [6–8]. Theoretical attempts to explain these data have been independently performed by several groups with the inclusion of the next-to-leading order QCD corrections [9–16].

Experimental and theoretical studies performed at hadron colliders have so far been limited to the $\Upsilon(1S, 2S, 3S)$ bound states, since they all lie below the $B\bar{B}$ threshold and hence have sizable leptonic branching fractions. Above the $B\bar{B}$ threshold, however, the leptonic branching ratios of the higher bottomonium states become very small, as a consequence of which these states have not been seen so far in hadronic collisions. But, if the anomalously large decay widths of $O(1)$ MeV in the final states $\Upsilon(1S, 2S, 3S)\pi^+\pi^-$, reported by Belle a few years ago [17, 18], are to be ascribed to the decays of the $\Upsilon(5S)$, then these final states are also promising for the detection of the $\Upsilon(5S)$ in experiments at the Tevatron and the LHC. Arguing along similar lines, the rescattering mechanism which enhances the dipionic partial widths in the $\Upsilon(5S)$ decays is also likely to yield similar enhancements in the rates for the corresponding transitions in the $\Upsilon(6S)$ decays [19], which then could also be measured in hadronic collisions. In this paper, we derive the hadroproduction cross sections for $\Upsilon(5S)$ and $\Upsilon(6S)$ in $p\bar{p}(p)$ collisions using the NRQCD framework, supplemented by the subsequent decays into $\Upsilon(1S, 2S, 3S)\pi^+\pi^-$.

The enhanced rates for the dipionic transitions in the $\Upsilon(5S)$-energy region could, however, also be ascribed to $Y_b(10890)$, a state reported by the Belle collaboration [17, 18], which is tentatively interpreted as a tetraquark [20–23]. In that case, one expects a smaller cross section for the hadroproduction of $Y_b(10890)$ than for a genuine $b\bar{b}$ bound state. At the same time, as there are no tetraquark states expected to lie in the $\Upsilon(6S)$ region, there would be no plausible grounds to expect a measurable yield in the $(\Upsilon(1S, 2S, 3S) \to \mu^+\mu^-)\pi^+\pi^-$ final states from the decays of $\Upsilon(6S)$. Since exotic states in the charm sector have been successfully searched for in the $(J/\psi, \psi')\pi^+\pi^-$ final states not only at the e^+e^- colliders, but also in hadroproduction in experiments at the Tevatron [24] and the LHC [25, 26], the proposed measurements at hadron colliders in the final states $(\Upsilon(1S, 2S, 3S) \to \mu^+\mu^-)\pi^+\pi^-$ could open new avenues in the search and discovery of the exotic four quark states in the bottom sector. In particular, there exist three candidates up to date, namely the states labeled $Y_b(10890)$, $Z_b(10610)$ and $Z_b(10650)$, with the last two observed by Belle last year [27]. If the exotic state $Y_b(10890)$ is not confirmed, we have nonetheless demonstrated a new way to explore the bottomonia above the $B\bar{B}$ threshold, which would supplement the study of $\Upsilon(1S, 2S, 3S)$ in hadronic collisions.

The cross section for the hadroproduction process $p\bar{p}(p) \to \Upsilon + X$ (we will leave X implicit in the following) is given by

$$\sigma_N(p\bar{p}(p) \to \Upsilon + X) = \int dx_1 dx_2 \sum_{i,j} f_i(x_1) f_j(x_2) \times \hat{\sigma}(ij \to \langle \bar{b}b\rangle_N + X)\langle O[N]\rangle, \quad (1)$$

where i, j denotes a generic parton inside a proton/antiproton, and $f_a(x_1), f_b(x_2)$ are the parton distribution functions (PDFs), which depend on the fractional momenta $x_i (i = 1, 2)$ (an additional scale-dependence is suppressed here), and Υ denotes a generic bottomonium state above $B\bar{B}$ threshold for which we consider $\Upsilon(5S)$ and $\Upsilon(6S)$ in this Letter. We adopt the CTEQ 6 PDFs [28] in our numerical calculations. $\langle O[N]\rangle$ are the long-distance matrix elements (LDMEs). N denotes all the quantum numbers of the $b\bar{b}$ pair, which we label in the form ${}^{2S+1}L_J^c$ (color c, spin S, angular momentum L, and total angular momentum J), and $\hat{\sigma}$ denotes the partonic cross section. The normalized cross sections, in which the LDMEs are factored out are defined by $\tilde{\sigma}_N \equiv \sigma_N/\langle O[N]\rangle$. The transverse momentum distribution is then given by

$$\frac{d\sigma_N}{dp_t} = \sum_{i,j} \int J dx_1 dy f_i(x_1, \mu_f) f_j(x_2, \mu_f) \frac{d\hat{\sigma}_N}{dt} \langle O[N]\rangle, \quad (2)$$

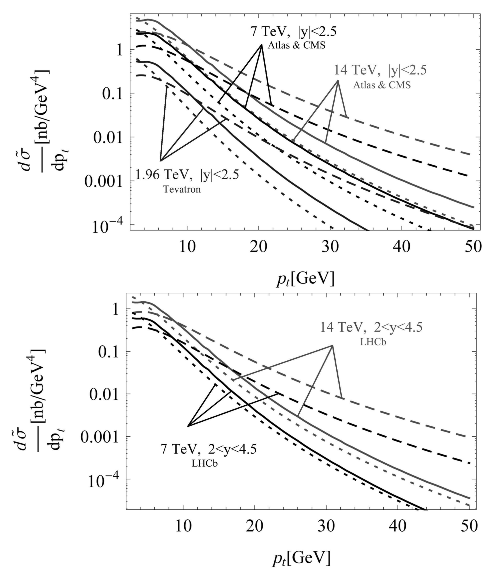

FIG. 1. Individual contributions (${}^3S_1^1$ solid, ${}^3S_1^8$ dashed, ${}^1S_0^8$ dotted, CO contributions are multiplied by 10^{-2}) for the normalized transverse momentum distributions $d\tilde{\sigma}/dp_t$ (explained in the text) for the process $p\bar{p}(p) \to \Upsilon(5S)$ (the corresponding curves for $\Upsilon(6S)$ are almost identical on a logarithmic plot, as are the distributions at 7 and 8 TeV). The p_t integrated values are given in Table I.

TABLE I. Integrated normalized cross sections $\tilde{\sigma}_N$, shown in Fig. 1 (in units of nb/GeV^3, CO channels are multiplied by 10^{-2}) for the processes $p\bar{p}(p) \to \Upsilon(5S, 6S)$, assuming a transverse momentum range 3 GeV $< p_t <$ 50 GeV. The rapidity range $|y| < 2.5$ has been assumed for the Tevatron experiments (CDF and D0) at 1.96 TeV and for the LHC experiments (ATLAS and CMS) at 7, 8 and 14 TeV; the rapidity range $2.0 < y < 4.5$ is used for the LHCb.

		$\Upsilon(5S)$			$\Upsilon(6S)$		
		${}^3S_1^1$	${}^3S_1^8$	${}^1S_0^8$	${}^3S_1^1$	${}^3S_1^8$	${}^1S_0^8$
Tevatron		2.72	1.73	1.75	2.54	1.66	1.60
LHC	7	13.25	9.49	9.06	12.44	9.16	8.32
LHCb	7	3.13	2.78	2.65	2.93	2.67	2.43
LHC	8	15.35	11.15	10.57	14.41	10.75	9.73
LHCb	8	3.80	3.35	3.17	3.56	3.22	2.92
LHC	14	27.62	21.15	18.76	25.98	20.48	17.30
LHCb	14	7.99	6.91	6.45	7.50	6.67	5.92

where y is the rapidity of Υ, p_t is the transverse momentum and J is the Jacobian factor.

The leading-order partonic processes for the S-wave configurations are:

$$\begin{aligned} g(p_1)g(p_2) &\to \Upsilon[{}^3S_1^1](p_3) + g(p_4) \\ g(p_1)g(p_2) &\to \Upsilon[{}^1S_0^8,\ {}^3S_1^8](p_3) + g(p_4), \\ g(p_1)q(p_2) &\to \Upsilon[{}^1S_0^8,\ {}^3S_1^8](p_3) + q(p_4), \\ q(p_1)\bar{q}(p_2) &\to \Upsilon[{}^1S_0^8,\ {}^3S_1^8](p_3) + g(p_4). \end{aligned} \quad (3)$$

These differential partonic cross sections, which are needed in Eq. (2) have been calculated in fixed-order perturbation theory in the literature. For the color singlet (CS), one has (see for instance Ref. [15]):

$$\frac{d\hat{\sigma}}{d\hat{t}} = \frac{5\pi^2\alpha_s^3[\hat{s}^2(\hat{s}-1)^2 + \hat{t}^2(\hat{t}-1)^2 + \hat{u}^2(\hat{u}-1)^2]}{216 m_b^5 \hat{s}^2(\hat{s}-1)^2(\hat{t}-1)^2(\hat{u}-1)^2}. \quad (4)$$

The Mandelstam variables are defined as

$$\hat{s} = \frac{(p_1+p_2)^2}{4m_b^2},\ \hat{t} = \frac{(p_1-p_3)^2}{4m_b^2},\ \hat{u} = \frac{(p_1-p_4)^2}{4m_b^2}, \quad (5)$$

with $m_b \simeq 4.75$ GeV. The factorization scale μ_f is chosen as $\mu_f = \sqrt{4m_b^2 + p_t^2}$. The partonic cross sections for the color octet (CO) have been calculated in Refs. [11, 29]. The K-factor for the CS contribution has been calculated for the process $p\bar{p}(p) \to \Upsilon(1S)$ in [15], which we have employed for the numerical calculations presented here. It is assumed, that the K-factor is not sensitive to $\sqrt{s}$. The CO contributions are taken at LO, since the NLO corrections, which have also been calculated for $p\bar{p}(p) \to \Upsilon(1S)$, are small [11].

Using these inputs, we show the transverse momentum distributions $\frac{d\sigma}{dp_t}$ in Fig. 1 for the processes $p\bar{p}(p) \to \Upsilon(5S)$ at the Tevatron (CDF and D0) and the LHC (ATLAS, CMS and LHCb) in the transverse momentum range 3 GeV $< p_t <$ 50 GeV, where $\log(p_t/m_{\Upsilon(5S)})$ is not

TABLE II. The obtained CO-LDMEs for bottomonia production (in unit of $10^{-2}\mathrm{GeV}^3$) from [16], extracted from data.

H	$< O^H\ ^1S_0^8 >$	$< O^H\ ^3S_1^8 >$
$\Upsilon(1S)$	11.15 ± 0.43	-0.41 ± 0.24
$\Upsilon(2S)$	3.55 ± 2.12	0.30 ± 0.78
$\Upsilon(3S)$	-1.07 ± 1.07	2.71 ± 0.13

large enough to necessitate the resummation of the logarithms [30–32]. The integrated normalized cross sections $\tilde{\sigma}_N$ are given in in Table I.

For the long-distance part we need nonperturbative input. The CS-LDMEs are given by the radial wave function at the origin and can be extracted from the partial e^+e^- widths via the Van-Royen Weisskopf formula. Using the Particle Data Group values [33] for the leptonic partial widths as input, and $m_{\Upsilon(5S)} = 10876$ MeV, $m_{\Upsilon(6S)} = 11019$ MeV, we find at NLO $|R(0)|^2_{\Upsilon(5S)} = 2.37\ \mathrm{GeV}^3$ and $|R(0)|^2_{\Upsilon(6S)} = 1.02\ \mathrm{GeV}^3$. The radial wave function at origin $R(0)$ is related to the LDME via $< O^H\ ^3S_1^1 >= 3|R(0)|^2/(4\pi)$.

The CO-LDMEs can only be extracted from the experimental data on differential distributions. This has been done for the $\Upsilon(1S, 2S, 3S)$ states by fitting the data on $\Upsilon(1S, 2S, 3S) \to \mu^+\mu^-$. The results for the CO matrix elements, extracted for $\Upsilon(1S, 2S, 3S)$ in Ref. [16], are displayed in Table II. We do not have the corresponding nonperturbative input for $\Upsilon(5S, 6S)$ at the current stage. Once the p_t-distributions in these states have been measured, the CO matrix elements can be extracted from an NRQCD-based analysis of the data. For our numerical estimates here, we will take the central values for the $\Upsilon(1S, 2S, 3S)$ states and assume as a first approximation, that the $\Upsilon(5S, 6S)$ have similar values. The spread in the values of the CO-LDMEs, given in Table II, is then taken as a rough error estimate.

For the exclusive production processes $p\bar{p}(p) \to \Upsilon(5S, 6S) \to (\Upsilon(1S, 2S, 3S) \to \mu^+\mu^-)\pi^+\pi^-$, we combine the results from the cross sections discussed earlier and the branching ratios for the $\Upsilon(nS)$ decays, which are listed in Table III. Note, that in this table, we have assumed that the Belle anomaly seen in the $\Upsilon(5S)$ decays can be explained in the rescattering model [19]. In this model, the $\Upsilon(6S) \to \Upsilon(1S, 2S, 3S)\pi^+\pi^-$ channels are also expected to have partial widths of about 1 MeV.

Using the available NRQCD results, we have explored the hadroproduction of bottomonium states above the $B\bar{B}$ threshold at the LHC and the Tevatron. The large branching fractions for the decays $\Upsilon(5S) \to \Upsilon(1S, 2S, 3S)\pi^+\pi^-$, observed by Belle [17, 18], offer an opportunity to access the $\Upsilon(5S)$ in hadronic collisions. Attributing the large enhancement in the dipionic transitions at the $\Upsilon(5S)$ to the rescattering phenomenon [19], very similar dipionic rates are expected for the $\Upsilon(6S)$ decays, which we have also worked out. In calculating the cross sections, we have included the next-to-leading order contributions by rescaling the available results for the process $p\bar{p}(p) \to \Upsilon(1S)$ [15]. The resulting cross sections presented here have large uncertainties, reflecting essentially our ignorance of the CO matrix elements for the $\Upsilon(5S)$ and $\Upsilon(6S)$, which we estimated using the LDMEs for $\Upsilon(1S, 2S, 3S)$ as input. However, once the experimental distributions for the $\Upsilon(5S)$ and $\Upsilon(6S)$ are available, the required CO matrix elements can be extracted from data, reducing the current uncertainties. Depending on the c.m. energy and the rapidity interval, our estimates yield typically a range of $O(10)$ pb, with a factor of 2 uncertainty up and down for the two LHC experiments ATLAS and the CMS. The cross sections are typically smaller by a factor 3 for the experiments at the Tevatron and LHCb. Nevertheless, given the current luminosities, they are large enough to undertake exploratory studies in hadronic collisions.

TABLE III. Branching ratios for the $\Upsilon(5S)$ and $\Upsilon(6S)$. All input values are taken from the PDG [33], except for the $\Upsilon(6S)$ entries, which are estimated from the scattering model [19].

$\mathcal{B}(\Upsilon(5S) \to \Upsilon(1S)\pi^+\pi^-)$	$(0.53 \pm 0.06)\%$
$\mathcal{B}(\Upsilon(5S) \to \Upsilon(2S)\pi^+\pi^-)$	$(0.78 \pm 0.13)\%$
$\mathcal{B}(\Upsilon(5S) \to \Upsilon(3S)\pi^+\pi^-)$	$(0.48 \pm 0.18)\%$
$\mathcal{B}(\Upsilon(6S) \to \Upsilon(1S)\pi^+\pi^-)$	$\approx 0.4\%$
$\mathcal{B}(\Upsilon(6S) \to \Upsilon(2S)\pi^+\pi^-)$	$(0.4 - 1.2)\%$
$\mathcal{B}(\Upsilon(6S) \to \Upsilon(3S)\pi^+\pi^-)$	$(1.2 - 2.5)\%$
$\mathcal{B}(\Upsilon(1S) \to \mu^+\mu^-)$	$(2.48 \pm 0.05)\%$
$\mathcal{B}(\Upsilon(2S) \to \mu^+\mu^-)$	$(1.93 \pm 0.17)\%$
$\mathcal{B}(\Upsilon(3S) \to \mu^+\mu^-)$	$(2.18 \pm 0.21)\%$

TABLE IV. Total cross sections for the processes $p\bar{p}(p) \to \Upsilon(5S, 6S) \to (\Upsilon(nS) \to \mu^+\mu^-)\pi^+\pi^-$ $(n = 1, 2, 3)$ in pb at the Tevatron ($\sqrt{s} = 1.96$ TeV) and the LHC ($\sqrt{s}$ =7, 8, 14 TeV), assuming the rapidity intervals described in Table I. The error estimates are from the variation of the central values of the CO-LDMEs, taken from $\Upsilon(1S, 2S, 3S)$, see text.

	$\Upsilon(5S)$			$\Upsilon(6S)$		
	$n = 1$	$n = 2$	$n = 3$	$n = 1$	$n = 2$	$n = 3$
Tevatron	2 ± 1	2 ± 1	1.2 ± 0.8	1.4 ± 1.0	2 ± 1	4 ± 3
LHC 7	9 ± 6	10 ± 7	7 ± 5	6 ± 4	9 ± 7	23 ± 17
LHCb 7	3 ± 2	3 ± 2	2 ± 1	2 ± 1	3 ± 2	7 ± 5
LHC 8	10 ± 7	12 ± 8	8 ± 5	7 ± 5	11 ± 8	27 ± 20
LHCb 8	3 ± 2	3 ± 2	2.4 ± 1.6	2.7 ± 2.0	3 ± 2	8 ± 6
LHC 14	19 ± 13	22 ± 15	15 ± 11	13 ± 10	20 ± 15	51 ± 39
LHCb 14	6 ± 4	7 ± 5	5 ± 3	4 ± 3	7 ± 5	17 ± 12

There are two competing scenarios, which can be explored in the future data analysis: i) Experiments are able to establish the signals in the processes $p\bar{p}(p) \to \Upsilon(5S, 6S) \to (\Upsilon(1S, 2S, 3S) \to \mu^+\mu^-)\pi^+\pi^-$, in agreement with the estimates presented here. ii) Experiments are able to establish only the process $p\bar{p}(p) \to \Upsilon(5S) \to (\Upsilon(1S,2S,3S) \to \mu^+\mu^-)\pi^+\pi^-$, but not $p\bar{p}(p) \to \Upsilon(6S) \to (\Upsilon(1S,2S,3S) \to \mu^+\mu^-)\pi^+\pi^-$, which, in our opinion, would speak against the rescattering mechanism and strengthen the case of $Y_b(10890)$ as the source of the anomalous dipion transitions [20–23].

Once enough data are available, one could undertake a Dalitz analysis of the $\Upsilon(1S,2S,3S)\pi^+\pi^-$ final states to determine the origin of the charged tetraquarks states $Z_b(10600)$ and $Z_b(10650)$, discovered by the BELLE collaboration [27] along similar lines. In this regard, we wish to point out that recently a charged four-quark state $Z_c(3900)$ has been discovered by the BESIII collaboration [34], confirmed by Belle [35], in the decays $Y(4260) \to Z_c(3900)^\pm\pi^\mp \to J/\psi\pi^+\pi^-$, where $Y(4260)$ is an exotic $c\bar{c}$ state [2], possibly a tetraquark [23, 36]. This lends indirect support to the interpretation that $Z_b(10600)$ and $Z_b(10650)$ are likewise the decay products of the exotic state $Y_b(10890)$. We note that $Z_c(3900)$ is also found in the analysis based on CLEO data [37]. These charmonium-like states can be accessed at hadron colliders in the final state $J/\psi^{(')}\pi^+\pi^-$.

In conclusion, we have shown by computing the cross sections for the processes $p\bar{p}(p) \to \Upsilon(5S,6S) \to (\Upsilon(1S,2S,3S) \to \mu^+\mu^-)\pi^+\pi^-$, that the experiments at the hadron colliders LHC and the Tevatron have the sensitivity to detect the bottomonium states $\Upsilon(5S)$ and $\Upsilon(6S)$, extending significantly their current experimental reach, and exploring thereby also the nature of the exotic states $Y_b(10890)$, $Z_b(10600)$ and $Z_b(10650)$, discovered in e^+e^- annihilation experiments.

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