Charmonium state formation and decay: $p\bar{p} \to {}^1D_2 \to {}^1P_1\gamma$

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Massless perturbative QCD forbids, at leading order, the exclusive annihilation of a proton-antiproton pair into some charmonium states, which, however, have been observed in the $p\bar{p}$ channel, indicating the significance of higher order and nonperturbative effects in the few GeV energy region. The most well known cases are those of the ${}^1S_0$ (η_c) and the ${}^1P_1$. The case of the ${}^1D_2$ is considered here and a way of detecting such a state through its typical angular distribution in the radiative decay ${}^1D_2 \to {}^1 P_1\gamma$ is suggested. Estimates of the branching ratio $B({}^1D_2 \to p\bar{p})$, as given by a quark-diquark model of the nucleon, mass corrections, and an instanton-induced process are presented.

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Several hadronic two-body decays of charmonium states are forbidden within leading-order perturbative QCD [1,2], but are nevertheless observed to occur with decay rates comparable to or even bigger than those of allowed decays [3]; the most well-known examples are the $J/\psi \to VP$ [4], and the $\eta_c \to B\bar{B}, VV$ [5] channels, where P is a pseudoscalar meson, V a vector meson, and B a baryon. Indeed the observed decay rates for $J/\psi \to \rho\pi, K^*\bar{K}$ and $\eta_c \to p\bar{p}, \rho\rho, \phi\phi, K^*\bar{K}^*$ are difficult to explain within conventional perturbative QCD (PQCD). Recently, also the ${}^1P_1$ coupling to $p\bar{p}$ has been established [6], despite being equally forbidden by the helicity conservation rule of massless PQCD [7].

Among the attempts to solve these problems, nonleading contributions [8], two quark correlations inside baryons [9], quark mass effects [10], and gluonic contents of the J/ψ [4] and the η_c [11] have been considered. Higher-order Fock states might help with the $J/\psi \to \rho\pi$ decay [8], but their contributions to other processes are not clear; diquarks and mass corrections do not help much with the η_c forbidden decays, whereas gluonic contributions seem to be more promising [12]. Recently a dynamical model for such contributions, with instanton-induced nonperturbative chiral symmetry breaking, has been used to obtain a good agreement with the data on $\Gamma(\eta_c \to p\bar{p})$ [13].

We consider here yet another case of a charmonium decay which should be forbidden according to PQCD, namely, ${}^1D_2 \to p\bar{p}$. Its observation would be very interesting because, among the above nonperturbative mechanisms invoked to explain the other forbidden decays, no one seems to be able to account for a sizable decay rate: as we shall see, both mass corrections and diquarks give very tiny decay rate values and instanton-induced processes are strongly suppressed with increasing Q^2 values [13].

Let us briefly recall why the ${}^1D_2 \to p\bar{p}$ decay is forbidden in massless PQCD. This charmonium state has quantum numbers $J^{PC} = 2^{-+}$: parity, angular momentum, and charge conjugation conservation only allow a final $p\bar{p}$ state with orbital angular momentum $L = 2$ and spin $S = 0$. $S = 0$ implies that the p and the $\bar{p}$ must have, in the charmonium rest frame, the same helicity, which is forbidden by the PQCD vector coupling of hard gluons to massless quarks and antiquarks. Such a helicity selection rule can only be broken by terms proportional to m_q/m_c or k_T/m_c, where m_q and m_c are, respectively, the light quark and charmed quark masses and k_T is the quark intrinsic transverse momentum. The current masses of u and d quarks are very small compared to the charmed quark mass and terms proportional to m_q/m_c are indeed negligible; terms proportional to k_T/m_c might be more relevant, but no comprehensive treatment of these contributions, together with other higher twist effects, has yet been performed.

Let us now consider the ${}^1D_2$ state created in $p\bar{p}$ annihilations, choosing the z axis as the proton direction in the $p\bar{p}$ center-of-mass frame. It is then clear, from what we said above, that the ${}^1D_2$ state can only be created with the spin third component $J_z = 0$; such charmonium state is then purely polarized and its spin-density matrix has only one nonzero component:

$$\rho_{00}({}^1D_2) = 1 \ . \tag{1}$$

This peculiar property reflects into the decay angular distributions of the ${}^1D_2$. One radiative decay which is expected to be observed with a large branching ratio is

$${}^1D_2 \to \ {}^1P_1\gamma \ , \tag{2}$$

which is dominated by an electric-dipole transition. The angular distribution of the photon, as it emerges in the rest frame of the ${}^1D_2$, is then simply given by [14]

$$W_\gamma(\theta) = \tfrac{1}{8}(5 - 3\cos^2\theta) \ , \tag{3}$$

where θ is the photon polar angle and an integration has been performed over the azimuthal angle.

The observation of such an angular distribution in $p\bar{p}$ exclusive annihilations should then be a clear signal of the formation and decay of the $^1D_2$ state; the full chain of processes to be looked for, according to the observed or expected decays of the $^1P_1$ state [6], is

$$p\bar{p} \to {}^1D_2 \to {}^1P_1\gamma \to (\eta_c\gamma)\gamma \to (\gamma\gamma\gamma)\gamma \ , \quad (4)$$

or

$$p\bar{p} \to {}^1D_2 \to {}^1P_1\gamma \to (J/\psi\pi^0)\gamma \to (e^+e^-\pi^0)\gamma \ . \quad (5)$$

Notice that the expected mass of the $^1D_2$ state is $M_D = (3788 \pm 7)$ MeV [15].

A most general analysis of the cascade processes (4) and (5) has been performed in Ref. [16]; in particular, from a study of the angular distribution of the two photons in the $^1D_2 \to {}^1P_1\gamma \to \eta_c\gamma\gamma$ decay one can obtain the values of the helicity amplitudes for the process $^1D_2 \to {}^1P_1\gamma$ and the expression of the photon angular distribution $W_\gamma(\theta)$. In case of dominance of the $E1$ electric-dipole transition (which one expects) the results of Ref. [16] agree with Eq. (3); in case other multipole amplitudes contribute (like $M2$ and $E3$) they give an explicit correction to Eq. (3).

So far the $^1S_0(\eta_c)$, $^3S_1(J/\psi$ and $\psi')$, $^3P_1(\chi_{c1})$, and $^3P_2(\chi_{c2})$ charmonium states have been observed to couple to $p\bar{p}$; the corresponding branching ratios $B(^{2S+1}L_J \to p\bar{p})$ are typically of the order of 10^{-4} to 10^{-3} [3]. Curiously, the $\eta_c \to p\bar{p}$ branching ratio, which should be zero according to lowest-order PQCD, is among the largest ones. Recently, also the $^1P_1$ has been observed in the $p\bar{p} \to {}^1P_1 \to J/\psi\pi^0$ channel [6], with an estimate for the product of the two branching ratios $B(^1P_1 \to p\bar{p})B(^1P_1 \to J/\psi\pi^0) \simeq 10^{-7}$. Notice that, similarly to what explained for the $^1D_2$, also the $^1P_1$ decay into $p\bar{p}$ is forbidden by leading order PQCD [7]. The $^3P_0$ state has not yet been observed, but this is presumably due to its small ($\lesssim 10^{-2}$) branching ratio into $J/\psi\gamma$; this makes the full process through which one looks for such a state, $p\bar{p} \to \chi_{c0} \to J/\psi\gamma \to e^+e^-\gamma$, a difficult one to detect. The analogous situation for the $^3P_1$ and $^3P_2$ states is much more favorable in that their branching ratios into $J/\psi\gamma$ are, respectively, $\simeq 0.27$ and 0.13 [3].

Thus, it is natural to expect a $^1D_2 \to p\bar{p}$ branching ratio similar to that observed for the other charmonium states. However, this would be very difficult to explain; to see why, we now briefly consider several possible nonperturbative contributions to such a process.

Mass corrections to "forbidden" charmonium decays have been considered in Ref. [10] for $\eta_c, \chi_{c0} \to p\bar{p}$: each quark inside the proton is assigned a constituent mass $m_q = xm_p$ and the c quarks in the charmonium state of mass M have a mass $m_c = M/2$. These corrections yield sizable values of $\Gamma(\chi_{c0} \to p\bar{p})$, but very small ones for $\Gamma(\eta_c \to p\bar{p})$, actually a factor $\sim 10^{-4}$ smaller than data. Following the same procedure and notations as in Ref. [10] we have computed the helicity amplitudes for the decay $^1D_2 \to p\bar{p}$; the only nonzero ones are

$$\begin{aligned}
A_{++;M}(\theta) &= -A_{--;M}(\theta) \\
&= \tfrac{32}{27}\sqrt{5/3}\,\pi^4\alpha_s^3 R''(0)\frac{F_N^2}{M_D^7}\epsilon(1-4\epsilon^2)^2 d^2_{M,0}(\theta)\int_0^1 dx_2\int_0^{1-x_2} dx_3\int_0^1 dy_2\int_0^{1-y_2} dy_3\frac{1}{x_2y_2+(x_2-y_2)^2\epsilon^2} \\
&\quad\times\frac{1}{1+x_2y_2-x_2-y_2+(x_2-y_2)^2\epsilon^2}\,\frac{1}{x_3y_3+(x_3-y_3)^2\epsilon^2}\,\frac{1}{(1-x_2)y_3+(1-x_2-y_3)^2\epsilon^2} \\
&\quad\times\frac{1}{[x_2y_2-\frac{1}{2}(x_2+y_2)+(x_2-y_2)^2\epsilon^2]^3}(x_2-y_2)^3 \\
&\quad\times\Bigg(-\Big\{\varphi_x(231)\varphi_y(321)-\varphi_x(132)[\varphi_y(321)+\varphi_y(123)]-[\varphi_x(132)+\varphi_x(231)]\varphi_y(123)\Big\}(1-x_2-y_3) \\
&\quad-\Big\{\varphi_x(123)\varphi_y(213)-\varphi_x(321)[\varphi_y(213)+\varphi_y(312)]-[\varphi_x(321)+\varphi_x(123)]\varphi_y(312)\Big\}(1-x_2) \\
&\quad+\Big\{\varphi_x(213)\varphi_y(123)-\varphi_x(312)[\varphi_y(321)+\varphi_y(123)]-[\varphi_x(213)+\varphi_x(312)]\varphi_y(321)\Big\}(1-x_2)\Bigg)\ , \quad (6)
\end{aligned}$$

where M_D is the $^1D_2$ mass and $R''(0)$ is the value at the origin of its wave-function second derivative. $\varphi_z(i,j,k) \equiv \varphi(z_i,z_j,z_k)$ denotes the proton distribution amplitude and F_N is a dimensional "decay constant" related to the value of the nucleon wave function at the origin; we refer to Ref. [10] for further details. Here we only notice that ϵ is the ratio of the proton to the charmonium mass, $\epsilon = m_p/M_D$, so that in the massless limit, $\epsilon \to 0$, indeed $A_{\pm\pm;M} = 0$, as required by PQCD.

From the knowledge of the decay helicity amplitudes one obtains the decay rate, upon introducing an explicit expression for the distribution amplitudes in Eq. (6) and performing the x and y integrations:

$$\begin{aligned}
\Gamma(^1D_2 \to p\bar{p}) &= \frac{(1-4\epsilon^2)^{1/2}}{40(2\pi)^4}\sum_{\lambda_p,\lambda_{\bar{p}},M}\int_{-1}^{1} d(\cos\theta)|A_{\lambda_p,\lambda_{\bar{p}},M}(\theta)|^2 \\
&= \frac{2^5}{3^7}\pi^4\alpha_s^6|R''(0)|^2|F_N|^4\epsilon^2(1-4\epsilon^2)^{9/2}\frac{I^2(\epsilon)}{M_D^{14}}\ , \quad (7)
\end{aligned}$$

where I is the multiple integral appearing in Eq. (6).

From[1] Ref. [9] also the decay rate of the $^1D_2$ into two gluons can be obtained:

$$\Gamma(^1D_2 \to gg) = \frac{32}{3}\frac{\alpha_s^2}{M_D^6}|R''(0)|^2 \; . \quad (8)$$

By assuming the total hadronic decay rate of the $^1D_2$ to be approximately given by Eq. (8), one can get an estimate of the branching ratio $B(^1D_2 \to p\bar{p})$ by taking the ratio of Eqs. (7) and (8), so that the unknown quantity $R''(0)$ cancels out. The result strongly depends on the choice of the distribution amplitudes $\varphi(x_1, x_2, x_3)$; according to the different choices adopted in Ref. [10] one obtains

$$B(^1D_2 \to p\bar{p}) \sim 10^{-8} - 10^{-12} \; . \quad (9)$$

Equation (9) clearly shows how mass corrections could not account for the eventual observation of the $^1D_2 \to p\bar{p}$ decay; the small values obtained for the branching ratio are mainly due to the factor $(x_2 - y_2)^3$ contained in the decay amplitude, Eq. (6). This is similar to what happens for the $\eta_c \to p\bar{p}$ process, where mass corrections are also very small, due to a factor $(x_2 - y_2)$ in the amplitude [10]; in the present case, actually, the situation is even worse, because of the third power of $(x_2 - y_2)$. In fact, in the η_c case, mass corrections lead to $B(\eta_c \to p\bar{p}) \sim 10^{-6}$–$10^{-10}$ [10], a result far away from the observed value $B(\eta_c \to p\bar{p}) \simeq 10^{-3}$, but bigger than the values given in Eq. (9).

One can similarly show that also two quark correlations could not explain a branching ratio for the $^1D_2 \to p\bar{p}$ decay of the order of 10^{-4}; a vector diquark component of the proton allows the decay, by allowing helicity flips at the gluon-vector diquark coupling [9], but, once more, the numerical values turn out to be too small. This can be explicitly checked by repeating the same procedure followed above for mass corrections; the expression of the decay helicity amplitudes, in the quark-diquark model of the proton, can be found in Ref. [9] and, again, it contains a small factor $(x - y)^3$. One finds, with little dependence on the choice of the distribution amplitudes,

$$B(^1D_2 \to p\bar{p}) \sim 10^{-8} \; . \quad (10)$$

Among other nonperturbative effects proposed to explain unexpectedly large branching ratios, the presence of the fundamental ($L = 0$) trigluonium states, with quantum numbers $J^{PC} = 0^{-+}, 1^{--}, 3^{--}$, in the 3-GeV mass region, has been proposed [4,11]. The first two states, mixing, respectively, with the η_c and the J/ψ, might explain some of their "mysterious" decays. However, a similar explanation for the $^1D_2$, the presence of a 2^{-+} glueball with a mass close to 3.8 GeV, looks much less natural and realistic.

Let us consider finally the instanton-induced mechanism proposed in Ref. [13] for the $\eta_c \to p\bar{p}$ decay: we know that its contribution decreases very rapidly with increasing Q^2 and, indeed, already for the decay of η_c', with a mass $\simeq 3.6$ GeV, is much smaller than for the η_c [13]. Considering the still higher mass of the $^1D_2$, $M_D \simeq 3.8$ GeV, we cannot expect this nonperturbative contribution to be large enough to produce a branching ratio for the process $^1D_2 \to p\bar{p}$ similar to those observed for the other charmonium states.

We have thus seen how several possible nonperturbative effects cannot contribute significantly to the $^1D_2$ coupling to $p\bar{p}$; on the other hand, we know that leading-order PQCD predicts no coupling at all, whereas higher-order corrections are difficult to evaluate and have never been computed. A similar situation occurs with the η_c, with the difference that for such particle one might expect a significant gluonic contribution [12,13]. Therefore, the eventual observation of a $B(^1D_2 \to p\bar{p}) \sim 10^{-4}$, analogous to the values observed for all other charmonium states which can couple to $p\bar{p}$, would pose an intriguing challenge to the theory.

The $^1D_2$ state could be looked for in the mass region $M_D \simeq 3788$ MeV [15] and in the reactions suggested by Eqs. (4) and (5), which should exhibit a typical decay angular distribution of the γ in the first step of the process. In fact the $^1D_2$ created in $p\bar{p}$ annihilation is in a pure spin state with $J_z = 0$ and its decay into $^1P_1\gamma$, dominated by an $E1$ transition, has the simple angular distribution given in Eq. (3). Actually, even if other multipole amplitudes contribute to this decay, their relative weights can be evaluated by looking at the angular distribution of the subsequent decay of the $^1P_1$ [16]. Hopefully, the $^1D_2 \to {}^1P_1\gamma$ radiative decay has a large branching ratio, so that the processes of Eqs. (4) and /or (5) can be detected. This is not unrealistic if one notices that the $^1D_2$ state, due to its expected mass and quantum numbers, cannot decay into pairs of D and/or D^* mesons.

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[1]Notice that in Ref. [9] the $^1D_2$ state is named f_2.

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