The decay $\eta_c \to p\bar{p}$ in a quark-diquark scheme

Mauro Anselmino, Francisco Caruso,* and Stefano Forte
Dipartimento di Fisica Teorica, Università di Torino, Torino, Italy
and Istituto Nazionale di Fisica Nucleare, Sezione di Torino, via P. Giuria 1, I-10125 Torino, Italy

Bernard Pire
Centre de Physique Theorique, Ecole Polytechnique, 91128 Palaiseau Cedex, France
(Received 17 February 1988)

The decay rate for $\eta_c \to p\bar{p}$ is computed at the tree level in perturbative QCD, modeling the proton with a quark-diquark system. Satisfactory agreement with the data is obtained. Comparisons with pure-quark models are made and predictions are given for this and other heavy-pseudoscalar-meson decay rates into baryon-antibaryon.

The correct description of exclusive processes in the framework of the quark-parton model and of perturbative QCD still presents many unsolved problems. This is essentially due to the fact that nonperturbative information has to be used in order to determine form factors and hadronization amplitudes into specific channels.

Diquarks have been recently suggested as an effective way of taking some of these effects into account. Indeed, whereas diquarks have been introduced a long time ago as static constituents in hadron spectroscopy,[1] more recently several arguments favoring a significant dynamical diquark substructure in baryons have been put forward both theoretically (as well in the context of hadron-hadron[2] as in deep-inelastic lepton-hadron scattering[3]) and experimentally.[4]

In particular, diquarks have been used as quasielementary constituents inside baryons in the description of p-p large-angle elastic scattering[2] and $p\bar{p}$ annihilation into hyperon-antihyperon,[5] leading to a good understanding of the experimental data. On the contrary, calculations inspired by pure quark QCD seem to run into trouble, especially when spin effects are involved. This is particularly true for exclusive reactions, since available data are not in the realm of asymptotic energies, where one expects to find a reasonable agreement between perturbative QCD predictions and the data. In the intermediate-energy region presently probed, diquarks do not seem to be resolved but act as elementary constituents.

In this paper we shall discuss the η_c decay into a baryon-antibaryon, and, in particular, the $\eta_c \to p\bar{p}$ decay, for which experimental data are available. This is the simplest process that presents a drastic spin effect, which, as we shall see, can be successfully described in a quark-diquark scheme, whereas in the pure massless quark model[6] it is forbidden.

While previous computations[2,5] have been carried out in the simplified end-point model (where one of the constituents carries most of the hadron momentum), here we shall use the full QCD scheme discussed in Ref. 6, adapted to the case of diquarks.

Since η_c is the lightest $c\bar{c}$ meson ($M_{\eta_c}=2980$ MeV, which excludes decays into particles with open charm) all its decay modes are suppressed by the Okubo-Zweig-Iizuka (OZI) rule. Although the available energy is not exceedingly large, these decays should be mediated by relatively hard gluons (otherwise, it would be difficult to explain the OZI suppression), thereby justifying a perturbative description in the QCD framework.

On the other hand, the energy scale for these processes is about the same as that of the available data on $\gamma\gamma \to p\bar{p}$, which appear to be rather well described by perturbative QCD with quark-diquark baryon substructure.[7] It seems therefore justified to model the exclusive decay modes of η_c with a ("hard") elementary process of the same kind.

The standard framework used to perform perturbative QCD calculations of exclusive reactions has been set and discussed in detail by Brodsky, Farrar, and Lepage (BFL) (Ref. 6). In this scheme, light-quark masses are neglected, and helicity is strictly conserved at each quark-gluon vertex.

Decays of the family $C \to p\bar{p}$, with C a charmonium state, have been described for some time with these methods.[8] Comparison with the experimental data[9] for the process $J/\psi \to p\bar{p}$ displays a fair agreement, although some (physically meaningful) parameters of the model may depend very strongly on the choice of wave function[10] (see below). Predictions for the branching ratio of $\chi_2 \to p\bar{p}$ (Ref. 11), however, are almost 2 orders of magnitude smaller than the observed values.

The most dramatic effect, however, appears in the decay $\eta_c \to p\bar{p}$: it is straightforward to see that in the BFL scheme this process is strictly forbidden. Indeed, spin and parity considerations force the final $p\bar{p}$ system to be in a total spin $S=0$ state, whereas helicity-conserving quark-gluon couplings can lead only to a final $p\bar{p}$ state with $S=1$. The decay, however, is observed experimentally with a branching ratio $B=(0.12\pm0.06)\%$.

Introducing vector diquarks as active constituents opens the possibility of helicity-flipping gluon-diquark couplings, thus allowing the decay. Of course, an alternative possibility would be that of using the pure-quark

BFL scheme with massive quarks, since this too allows helicity-flipping quark-gluon couplings, proportional to the light-quark mass, a method already used to improve the computation of $\Gamma(J/\psi \to p\bar{p})$ (Ref. 12). This, however, leads to a very poor agreement with the data, and to some inconsistencies, as we shall discuss later. In fact, it is our aim to test the successfulness of our diquark scheme against that of pure quark QCD, even allowing for massive quarks.

Within the BFL scheme adapted to diquarks, the helicity amplitudes for $\eta_c \to p\bar{p}$ can be written as

$$M_{\lambda_p,\lambda_{\bar{p}}}(\eta_c \to p\bar{p}) = \sum_{\text{hel,col,fl}} \int dx\, dy\, dz\, \psi^*_{p,\lambda_p}(x)\psi^*_{\bar{p},\lambda_{\bar{p}}}(y) T_{H\lambda_q,\lambda_Q,\lambda_{\bar{q}},\lambda_{\bar{Q}};\lambda_c,\lambda_{\bar{c}}}(c\bar{c} \to qQ\bar{q}\bar{Q};x,y,z)\psi_{\eta_c}(z) \;, \tag{1}$$

where the ψ's are the flavor, color, spin, and space wave functions of the hadrons and T_H is the elementary helicity amplitude for the constituent process $c\bar{c} \to qQ\bar{q}\bar{Q}$ (in our scheme where a nucleon is made up of a quark q and a diquark Q). The sum goes over all allowed quantum numbers of the constituents (colors, helicities and flavors). The Feynman diagrams corresponding to T_H are drawn on Fig. 1, where the kinematics is also defined (in the rest frame of η_c). The η_c wave function is taken as the nonrelativistic one

$$\psi_{\eta_c}(z) = \frac{F_{\eta_c}}{\sqrt{2}}(c_+\bar{c}_- - c_-\bar{c}_+)\delta(z - \tfrac{1}{2}) \;, \tag{2}$$

where the momentum of η_c is equally divided between the heavy c and $\bar{c}$ quarks, and we have omitted the usual color part.

The quark-diquark proton wave function is given by[7]

$$\psi_{p,\lambda_p=\pm 1/2}(x) = \frac{\pm F_N}{\sqrt{18}}\{\phi_2(x)[\sqrt{2}V_{\pm 1}(ud)u_{\mp} - 2V_{\pm 1}(uu)d_{\mp}] + \phi_3(x)[\sqrt{2}V_0(uu)d_{\pm} - V_0(ud)u_{\pm}]$$
$$\mp [2\phi_1(x) + \phi_3(x)]S(ud)u_{\pm}\} \;, \tag{3}$$

where $V_h(ud)$ stands for a vector diquark made of a u and d quark with helicity h and so on (S = scalar diquark). F_{η_c} and F_N are unknown constants with the dimensions of (mass), whose squared moduli are somehow related to the probability for the quark-quark or quark-diquark pair to hadronize into the pertinent meson or baryon; they are the analogue of the pion decay constant F_π. We have written ψ_p assuming, for simplicity, its different components (scalar and vector diquarks with helicity ± 1 and 0) to have the same values of F_N. This is only approximately true since the strict equality is violated by several effects, as, for instance, mass differences between scalar and vector diquarks. The derivation of the wave function (3) is described in Ref. 7 and is based on the ansatz of Refs. 13 and 14. It reduces to the usual SU(6) wave function when $\phi_1 = \phi_2 = \phi_3$.

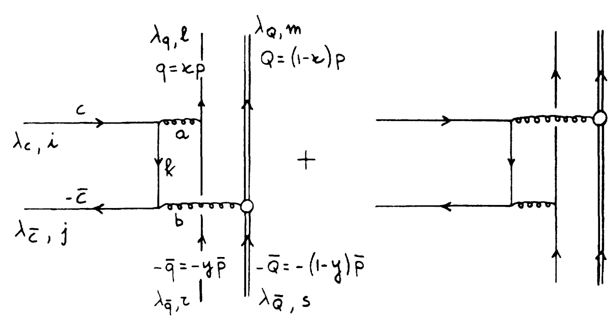

FIG. 1. Feynman diagrams for the elementary process $c\bar{c} \to qQ\bar{q}\bar{Q}$. Here

$$p^\mu = (M, p\sin\theta, 0, p\cos\theta), \quad \bar{p}^\mu = (M; -p\sin\theta, 0, -p\cos\theta) \;,$$

$$c^\mu = \bar{c}^\mu = (M,0) = \tfrac{1}{2}(M_{\eta_c}, 0), \quad \text{and } M_{\eta_c} = 2.980 \text{ GeV} \;.$$

i,j,l,m,r,s,a,b are color indices. $\lambda_q, \lambda_{\bar{q}}, \lambda_Q, \lambda_{\bar{Q}}$ are helicities and $\lambda_c, \lambda_{\bar{c}}$ are the z components of $S_c, S_{\bar{c}}$, respectively.

Throughout our calculation (see Fig. 1), we take parton masses into account and neglect the Fermi motion of the constituents. We also have to assign the diquark a definite mass, which appears explicitly in the diquark polarization vectors. The only consistent way of doing so is to use the naive parton model, that is, to assign to quarks and diquarks a running mass $m_q = xm_p$ and $M_Q = (1-x)m_p$, respectively (see, e.g., Ref. 15). We neglect the Q^2 dependence of the wave functions due to QCD evolution, for we are considering only a very limited energy range (up to the mass of the η_c).

In order to compute the elementary amplitudes we need to know the coupling of diquarks to gluons. These are given, in the most general (gauge-invariant) form, in Fig. 2, both for scalar and vector diquarks. F_s and G_i ($i = 1,2,3$) are the diquark form factors; the pointlike couplings are recovered in the small-Q^2 limit: for a scalar diquark

$$F_s(Q^2) \underset{Q^2 \to 0}{\longrightarrow} 1$$

and for a vector diquark with anomalous color-magnetic moment k, $G_1(Q^2) \to 1$, $G_2(Q^2) \to 1 + k$, and $G_3(Q^2) \to 0$. The only nonzero elementary amplitudes turn out to be

$$
\begin{aligned}
&T_{H\lambda_q,\lambda_Q\neq 0,\lambda_{\bar q},\lambda_{\bar Q}=0;\lambda_c,\lambda_{\bar c}}(x,y)=-8i\sqrt{2}g_s^4M^3p^2(2\lambda_c)(\lambda_Q-2\lambda_q)\delta_{\lambda_c,-\lambda_{\bar c}}\delta_{\lambda_q,-\lambda_{\bar q}}\frac{(1-y)(y-x)}{m_pg_1^2g_2^2(k^2-M^2)}C_FG_2(g_2^2)\ ,\\
&T_{H\lambda_q,\lambda_Q=0,\lambda_{\bar q},\lambda_{\bar Q}\neq 0;\lambda_c,\lambda_{\bar c}}(x,y)=8i\sqrt{2}g_s^4M^3p^2(2\lambda_c)(\lambda_{\bar Q}+2\lambda_q)\delta_{\lambda_c,-\lambda_{\bar c}}\delta_{\lambda_q,-\lambda_{\bar q}}\frac{(1-x)(y-x)}{m_pg_1^2g_2^2(k^2-M^2)}C_FG_2(g_2^2)\ ,
\end{aligned}
\tag{4}
$$

where

$$
\begin{aligned}
g_1^2&=(x-y)^2m_p^2+4xyM^2\ ,\\
g_2^2&=(x-y)^2m_p^2+4(1-x)(1-y)M^2\ ,\\
k^2-M^2&=(x-y)^2m_p^2+2(2xy-x-y)M^2\ ,
\end{aligned}
\tag{5}
$$

m_p and M are, respectively, the proton and charmed-quark masses, and C_F is a color factor (which already takes into account the color wave functions of η_c, p, and $\bar p$), $C_F=2\sqrt{3}/9$.

It should be noted that all elementary amplitudes involving scalar diquarks are zero regardless of whether or not quark and diquark running masses are taken into account. The process thus goes exclusively through the vector diquark coupling proportional to the form factor G_2 (see Fig. 2). By inserting Eqs. (4) into Eq. (1), and using Eqs. (2) and (3) we get

$$
\begin{aligned}
M_{++}(\eta_c\to p\bar p)&=-M_{--}\\
&=-\frac{i\pi^2 2^{10}\sqrt{6}}{m_p 3^3}F_{\eta_c}F_N^2M^3(M^2-m_p^2)I\ ,\\
M_{+-}&=M_{-+}=0\ ,
\end{aligned}
\tag{6}
$$

where

$$
I=\int dx\,dy\frac{\phi_2(x)\phi_3(y)\alpha_s^2(1-x)(y-x)}{g_1^2g_2^2(k^2-M^2)}G_2(g_2^2)\ . \tag{7}
$$

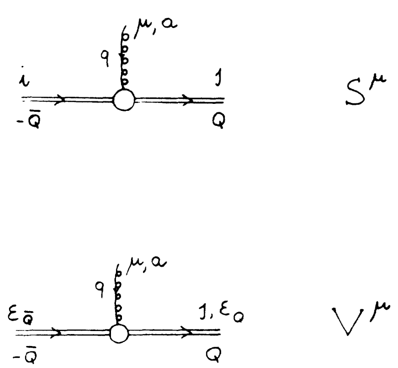

FIG. 2. The most general gluon couplings to scalar and vector diquarks. T^a are color Gell-Mann matrices; $\epsilon_Q,\epsilon_{\bar Q}$ are diquark polarization vectors. F_s, G_1, G_2, and G_3 are form factors:

$$
\begin{aligned}
S^\mu&\equiv -ig_sT_{ij}^a(Q-\bar Q)^\mu F_s(q^2)\ ,\\
V^\mu&\equiv ig_sT_{ij}^a\{(\epsilon_Q^*\cdot\epsilon_{\bar Q}^*)(Q-\bar Q)^\mu G_1(q^2)\\
&\quad-[(Q\cdot\epsilon_{\bar Q}^*)\epsilon_Q^\mu-(\bar Q\cdot\epsilon_Q^*)\epsilon_{\bar Q}^\mu]G_2(q^2)\\
&\quad-(\epsilon_Q^*\cdot\bar Q)(\epsilon_{\bar Q}^*\cdot Q)(Q-\bar Q)^\mu G_3(q^2)\}\ .
\end{aligned}
$$

Finally, from

$$
\Gamma(\eta_c\to p\bar p)=\frac{(M^2-m_p^2)^{1/2}}{32\pi M^2}\sum_{\lambda_p,\lambda_{\bar p}}|M_{\lambda_p,\lambda_{\bar p}}(\eta_c\to p\bar p)|^2 \tag{8}
$$

and Eq. (6) we have

$$
\Gamma(\eta_c\to p\bar p)=\frac{2^{17}\pi^3M^4}{3^5m_p^2}(M^2-m_p^2)^{5/2}|F_{\eta_c}|^2|F_N|^4I^2\ . \tag{9}
$$

This yields directly the decay rate, at the lowest perturbative level, in terms of the hadronization constants F_{η_c} and F_N, defined in Eqs. (2) and (3). These can be fixed by comparing with the data the theoretical predictions for different processes involving the same hadrons.

We have fixed F_{η_c} in two independent ways. First, we considered the electromagnetic decay $\eta_c\to\gamma\gamma$. Its amplitude is obtained by convoluting the elementary amplitudes corresponding to the Feynman diagrams of Fig. 3 with the η_c wave function (2), and is

$$
M_{\lambda_1\lambda_2}(\eta_c\to\gamma\gamma)=\frac{32\sqrt{3}}{9}i\pi\alpha\lambda_1\delta_{\lambda_1,\lambda_2}F_{\eta_c}\ . \tag{10}
$$

This gives

$$
\Gamma(\eta_c\to\gamma\gamma)=\frac{64\pi\alpha^2}{27M}|F_{\eta_c}|^2\ . \tag{11}
$$

Comparing with the experimental value[16]

$$
\Gamma(\eta_c\to\gamma\gamma)=5.7\pm 6.3\ \text{keV} \tag{12}
$$

yields

$$
|F_{\eta_c}|=146\pm 81\ \text{MeV}\ . \tag{13}
$$

Alternatively, we have fixed $|F_{\eta_c}|$ by assuming $\Gamma(\eta_c\to\text{all})\simeq\Gamma(\eta_c\to gg)$ with

$$
\Gamma(\eta_c\to gg)=\frac{16\pi\alpha_s^2}{27M}|F_{\eta_c}|^2\ , \tag{14}
$$

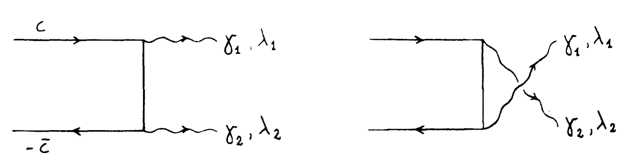

FIG. 3. Feynman diagrams for the elementary process $c\bar c\to\gamma\gamma$. λ_1,λ_2 are helicities.

where the color factor is the same as that of $\eta_c \to p\bar{p}$, since the two gluons hadronize into color-singlet states. From the experimental total width $\Gamma_{\eta_c} = 11 \pm 4$ MeV we get [in fair agreement with (13)]

$$|F_{\eta_c}| \leq 335 \pm 61 \text{ MeV} . \quad (13')$$

Note that, since *a priori* contributions from higher-order processes to $\Gamma(\eta_c \to \text{all})$ cannot be excluded, this should be understood as a lower limit on Γ, i.e., as an upper bound on $|F_{\eta_c}|$.

$|F_N|$ has been fixed by computing the electromagnetic form factor of the proton and the cross section for the process $\gamma\gamma \to p\bar{p}$ within a quark-diquark model for the proton (see Ref. 7 for a detailed discussion). This yields the values

$$|F_N| = 80\text{–}110 \text{ MeV} \quad (15)$$

according to the different wave functions used. The computation of Ref. 7 is performed neglecting the contribution of vector diquarks, which is estimated to be nonleading. Since the experimental errors and the statistical uncertainty of the fit are quite large, this determination of F_N is affected by a significant error.

On the other hand, both the $\gamma\gamma \to p\bar{p}$ cross section and the $\eta_c \to p\bar{p}$ decay rate are rather sensitive to the choice of F_N (they are proportional to $|F_N|^4$). Moreover the assumption of a unique overall value of F_N in (3) is an approximation. As a consequence, rather than choosing quite arbitrarily a value for F_N and getting a prediction for the η_c decay rate with huge error bars, we prefer to express our results as a determination of F_N. This yields a prediction (with a relatively small error) which can be compared with the available value of F_N as a first consistency check of the model. As soon as more data will become available, both theoretically and experimentally, it will be possible to turn the consistency check into a set of predictions.

In order to compute the decay amplitude for $\eta_c \to p\bar{p}$ we still have to choose a quark-diquark space wave function for the proton. A first simple possibility is to choose wave functions of the general form

$$\phi_2(x) = \phi_3(x) = kx^{\alpha}(1-x)^{\beta} \quad (16)$$

including, in particular,[7] (a) asymptotic wave function

$$\phi_2(x) = \phi_3(x) = 20x(1-x)^3 \quad (17a)$$

and (b) symmetric ("mesonlike") wave function

$$\phi_2(x) = \phi_3(x) = 6x(1-x) . \quad (17b)$$

A third choice, consistent with QCD sum rules, is (c) the Chernyak-Zhitnitsky-type wave function

$$\phi_2(x) = 20x(1-x)^3(5.292 - 24.864x + 27.972x^2) ,$$
$$\phi_3(x) = 20x(1-x)^3(5.544 - 27.888x + 33.264x^2) . \quad (17c)$$

[Equation (17c) can be derived following the procedure described in the Appendix of Ref. 7 and choosing for $\varphi(x_1,x_2,x_3)$ the expressions given in Ref. 14.]

The diquark form factor G_2 is, in our Q^2 region, of the form[1,2] $(1+k)[Q_0^2/(Q_0^2+Q^2)]$. The forward normalization is expressed in terms of the anomalous magnetic moment of the diquark k. In a naive additive quark model $k \approx 1$ (Ref. 17). Since in our energy scale $\langle Q^2 \rangle \simeq Q_0^2 \simeq 3$ GeV^2, it follows that $\langle G_2 \rangle \approx 1$ is a good approximation, as we have verified by explicit calculation using the wave functions (17a)–(17c); this supports the idea that, at the moderate energy range we are considering here, the diquarks behave as quasielementary objects.

Finally, we take a fixed value of α_s: namely,

$$\alpha_s(m_{\eta_c}^2) = \frac{12\pi}{25\ln(m_{\eta_c}^2/\Lambda^2)} \simeq 0.28 \quad (\Lambda = 0.2 \text{ GeV}) .$$

Comparing the decay rate (9) with the experimental value

$$\Gamma(\eta_c \to p\bar{p}) = (1.21 \pm 0.79) \times 10^{-2} \text{ MeV}$$

we find, using the wave functions (a), (b), and (c) and the value (13) of $|F_{\eta_c}|$,

$$|F_N| = 330 \pm 105 \text{ MeV} , \quad (18a)$$

$$|F_N| = 297 \pm 95 \text{ MeV} , \quad (18b)$$

$$|F_N| = 249 \pm 80 \text{ MeV} . \quad (18c)$$

Using instead the value (13′) of F_{η_c} one obtains

$$|F_N| = 218 \pm 29 \text{ MeV} , \quad (18a')$$

$$|F_N| = 196 \pm 27 \text{ MeV} , \quad (18b')$$

$$|F_N| = 164 \pm 23 \text{ MeV} , \quad (18c')$$

which should be taken as lower limits for $|F_N|$. These results have several interesting features. First, it is apparent that the value of $|F_N|$ is rather insensitive to the choice of wave function. Indeed, we have repeated the computation of the width with wave functions of the form (16) varying α and β from 0.1 to 10. We have verified that the amplitudes change at most by a factor 3; i.e., F_N varies by a factor $\sqrt{3}$. Note that with $0 < (\alpha,\beta) < 1$ the amplitude displays end-point divergences (although the integrated width is finite); but this, too, does not affect the result.

Our determination (18) and (18′) of F_N are consistent with each other. They are, with large errors, systematically higher, by a factor 2–3, than the determinations of F_N (Ref. 4) given in Eq. (15). As we already said, however, only vector diquarks contribute to $\eta_c \to p\bar{p}$, whereas only scalar ones have been taken into account in deriving the values (15), and there is no need why the hadronization constants should be exactly the same in the two cases.

Actually, we have repeated the computation by introducing a running coupling constant α_s [$\alpha_s(g_1^2)$ and $\alpha_s(g_2^2)$ with a maximum allowed valued of $\alpha_s = 0.5$] as done in Refs. 7 and 10. This leads to values of $|F_N|$ systematically smaller by roughly a factor $\sqrt{2}$. The determination of $|F_N|$, Eq. (15), was made with this choice of α_s, and indeed much better agreement is found when the

TABLE I. Predictions for the ratio $R_B \equiv \Gamma(\eta_c \to B\bar{B})/\Gamma(\eta_c \to p\bar{p})$ in a quark-diquark scheme. The first three columns of R_B are calculated assuming $F_B \simeq F_p$. The last three assume F_B/F_p to scale like m_B/m_p. (a)–(c) refer to the different wave functions, defined in Eq. (17) of the text.

		$\Gamma(\eta_c \to B\bar{B})/\Gamma(\eta_c \to p\bar{p})$					
B	m_B (MeV)	(a)	(b)	(c)	(a)	(b)	(c)
Λ	1115.60	0.27	0.26	0.23	0.54	0.52	0.46
Σ^+	1189.37	0.14	0.13	0.11	0.36	0.34	0.28
Σ^0	1192.46	0.14	0.13	0.11	0.37	0.34	0.29
Σ^-	1197.34	0.12	0.12	0.11	0.32	0.32	0.27
Ξ^0	1314.90	0.03	0.03	0.02	0.11	0.11	0.08
Ξ^-	1321.32	0.03	0.02	0.02	0.12	0.08	0.08

values given in Eq. (18) are divided by $\sqrt{2}$ (although we believe a fixed α_s to be more appropriate to this case). Furthermore, in Ref. 7, in order to compare with Ref. 10, we have neglected all masses (except the diquark ones); the same kind of approximation here would lead to yet significantly smaller values of F_N.

Let us now see how our results compare to a pure-quark picture, with massive quarks. In this scheme the only helicity amplitudes which contribute to the process are those where two quarks connected by a gluon replace the diquark in the diagrams of Fig. 1, with a total helicity of the quark-quark pair equal to the vector diquark polarizations of the nonzero amplitudes (4).

The resulting decay rate, however, is suppressed by an additional factor of α_s^2 (the process is a higher-order one) and by $(m_q/E)^2 \sim (2m_p/3M_{\eta_c})^2$ (the helicity flip is a mass effect), thus, overall by a factor ~ 0.05. Moreover, the value of F_N in a quark scheme, as derived from QCD sum rules, should be[13,14] $|F_N^q| = 5.2\times 10^{-3}$ GeV^2. One would expect F_N^q to be roughly connected to the quark-diquark F_N by $|F_N^q| \simeq |F_N|\,|F_D|$, where F_D is the "decay constant" of the diquark into two quarks. A simple estimate $|F_D| \sim 40$ MeV (as for ρ mesons[6]) implies, from the above value of $|F_N^q|$, $F_N \sim 130$ MeV in qualitative agreement with our results. It follows that the prediction from a pure quark scheme is suppressed by the above factor 0.05 as compared to our result.

It is interesting to remark that the phenomenological values of F_N, used to describe $\gamma\gamma \to p\bar{p}$ in pure quark models,[10] which are based on the computation of the ratio $\Gamma(J/\psi \to p\bar{p})/\Gamma(J/\psi \to e^+e^-)$, are very strongly dependent on the choice of wave function (varying by up to 5 orders of magnitude[10]), and are based on fits to the angular distribution of $J/\psi \to p\bar{p}$ which are in very dubious agreement with the data unless quark masses are taken into account.[12,9]

To convince oneself that the difference between the two schemes is quite independent of the definition of the hadronization constant, and is rather a consequence of the fundamental description of the elementary process, one may consider the ratio $\Gamma(J/\psi \to \gamma p\bar{p})/\Gamma(J/\psi \to p\bar{p})$, which does not depend on any F_N-like constant. A simple counting rule shows that in the pure quark picture this ratio is of the order of $\alpha = \frac{1}{137}$, while in a quark-diquark scheme it goes as $\alpha/\alpha_s^2 \simeq 0.1$. Experimentally, it is found to have the value[16(b)] $\simeq 0.2$.

We have thus shown that the quark-diquark scheme consistently describes the $\eta_c \to p\bar{p}$ decay rate. This gives additional evidence that when spin effects are involved, vector diquarks provide the helicity-flip amplitudes which are required to account for the experimental results.

Now, assuming SU(6)-like quark-diquark wave functions for baryons, we are able to predict the η_c decay rates into baryon-antibaryon as the ratio

$$R_B \equiv \frac{\Gamma(\eta_c \to B\bar{B})}{\Gamma(\eta_c \to p\bar{p})} = \left[\frac{m_p}{m_B}\right]^2 \left[\frac{M^2 - m_B^2}{M^2 - m_p^2}\right]^{5/2} \frac{|F_B|^4}{|F_p|^4}\frac{I^2(m_B)}{I^2(m_p)}, \tag{19}$$

where $I(m)$ is given by Eq. (7). The numerical results are shown in Table I and R_B is found to be actually insensitive to the choice of the proton wave function.

As a last application of our model, we compute the $\eta_b(b\bar{b})$ decay rates into $B\bar{B}$ in the form of the ratio $\Gamma(\eta_b \to B\bar{B})/\Gamma(\eta_c \to B\bar{B})$. It is assumed throughout that $|F_{\eta_b}|/|F_{\eta_c}|$ scales as M_b/M_c. Results are given in Table II and show a very interesting feature: in contrast with R_B, they strongly depend on the choice of the wave functions. This implies that precise measurements of two different particle (η_c, η_b) decays into the same final state $(B\bar{B})$ will give us the possibility of a comparative study of the quark-diquark hadronic wave functions.

We can safely conclude that our model is consistent with the current data. Our results are affected by several uncertainties, characteristic of perturbative QCD computations of exclusive processes: the choice of coupling constant, the value of the hadronization constants, and the functional form of the hadronic wave function. Addi-

TABLE II. Prediction for a pseudoscalar $b\bar{b}$ meson (η_b) decay into $B\bar{B}$.

		$\Gamma(\eta_b \to B\bar{B})/\Gamma(\eta_c \to B\bar{B})$	
Baryon	(a)	(b)	(c)
p	0.021	0.203	0.076
Λ	0.052	0.495	0.187
Σ^+	0.087	0.823	0.310
Σ^0	0.090	0.844	0.318
Σ^-	0.093	0.878	0.330
Ξ^0	0.314	2.944	1.097
Ξ^-	0.344	3.220	1.198

tional sources of uncertainty, peculiar to our treatment, are the choice of diquark form factor, and the use of running quark and diquark masses. Note, however, that we have argued that the form factor is approximately constant and equal to one in the energy range we are interested in.

We hope that new experimental information, allowing a better determination of the phenomenological parameters and a stricter comparison with the full set of our predictions will become available soon. Furthermore, the computation of more processes within our quark-diquark scheme would allow us to turn our qualitative results into a set of firm quantitative predictions. In particular, many more decays of the charmonium, bottomonium, and toponium family can be described in the quark-diquark model. Work is in progress towards a systematic treatment of heavy-quarkonium decays within this approach.

ACKNOWLEDGMENTS

We would like to thank E. Leader, D. B. Lichtenberg, E. Maina, and E. Predazzi for many discussions. M.A. and F.C. would like to thank the Theory Group of Ecole Polytechnique and the members of the Dipartimento di Fisica Teorica of the University of Torino, respectively, for their kind hospitality. F.C. also thanks the Conselho Nacional de Pesquisas do Brazil for financial support.

---

[*]On leave of absence from Institute of Physics, Universidade do Estado do Rio de Janeiro, Brazil.

[1]M. Ida and R. Kobayashi, Prog. Theor. Phys. **36**, 846 (1966); D. B. Lichtenberg and L. J. Tassie, Phys. Rev. **155**, 1601 (1967); D. B. Lichtenberg, W. Namgung, E. Predazzi, and J. G. Wills, Phys. Rev. Lett. **48**, 1653 (1982); A. Martin, Z. Phys. C **32**, 359 (1986).

[2]See, e.g., M. Anselmino, P. Kroll, and B. Pire, Z. Phys. C **36**, 86 (1987), and references therein.

[3]See, e.g., L. F. Abbott, E. L. Berger, R. Blankenbecler, and G. L. Kane, Phys. Lett. **88B**, 157 (1979), and references therein.

[4]A. Breakstone *et al.*, Z. Phys. C **28**, 335 (1985); **36**, 567 (1987); J. L. Bailly *et al.*, *ibid.* **31**, 367 (1986); **35**, 295 (1987); C. C. Chang *et al.*, Phys. Rev. D **27**, 2776 (1983); M. Barth *et al.*, Z. Phys. C **37**, 187 (1981).

[5]P. Kroll and W. Schweiger, Nucl. Phys. **A474**, 608 (1987).

[6]S. J. Brodsky and G. R. Farrar, Phys. Rev. D **11**, 1309 (1975); G. P. Lepage and S. J. Brodsky, *ibid.* **22**, 2157 (1980).

[7]M. Anselmino, F. Caruso, P. Kroll, and W. Schweiger, CERN Report No. TH-4941, 1988 (unpublished).

[8]S. J. Brodsky and G. P. Lepage, Phys. Rev. D **24**, 2848 (1981).

[9]D. Pallin *et al.*, Nucl. Phys. **B292**, 653 (1987).

[10]G. R. Farrar, E. Maina, and F. Neri, Nucl. Phys. **B259**, 702 (1985).

[11]P. Damgaard, K. Tsokos, and E. L. Berger, Nucl. Phys. **B259**, 285 (1985).

[12]C. Carimalo, Int. J. Mod. Phys. **A2**, 249 (1987).

[13]V. L. Chernyak and I. R. Zhitnitsky, Nucl. Phys. **B246**, 52 (1984); Phys. Rep. **112**, 173 (1984).

[14]I. D. King and C. T. Sachrajda, Nucl. Phys. **B279**, 785 (1987).

[15]E. Leader and E. Predazzi, *An Introduction to Gauge Theories and the New Physics* (Cambridge University Press, Cambridge, England, 1984), Chaps. 13 and 14.

[16](a) R. M. Baltrusaitis *et al.*, Phys. Lett. **106B**, 497 (1981); (b) Particle Data Group, M. Aguilar-Benitez *et al.*, *ibid.* **170B**, 1 (1986); (c) C. Baglin *et al.*, Phys. Lett. B **187**, 191 (1987).

[17]D. B. Lichtenberg (private communication).