Mass corrections in $J/\psi \to B\bar{B}$ decay and the role of distribution amplitudes

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(Received 30 November 1994)

We consider constituent quark mass corrections to the polar angular distribution of a baryon-antibaryon pair created in the chain decay process $e^- e^+ \to J/\psi \to B\bar{B}$, generalizing a previous analysis of Carimalo. We show the relevance of the features of the baryon distribution amplitudes and estimate the electromagnetic corrections to the QCD results.

PACS number(s): 13.25.Gv, 12.38.Bx, 13.65.+i, 14.40.Gx

I. INTRODUCTION

In the last few years our understanding of exclusive hadronic processes at high momentum transfer, in the framework of perturbative QCD (PQCD), has improved (for a comprehensive review and further references see, e.g., [1]). Theoretical models, based essentially on factorization ideas, have been elaborated and refined. The main ingredients of these models (which from now on we shall indicate as PQCD models) can be summarized as follows. The amplitude for a given exclusive process is obtained by convoluting two well distinct contributions: the first coming from the hard scattering among the partonic constituents of the involved hadrons and the second from the subsequent, soft processes which lead to hadronization. The hard scattering can be described by means of perturbative QCD techniques, representing, in first approximation, each participating hadron by its valence constituents, assumed collinear with the parent hadron and among themselves. As for the soft processes, their treatment is outside the possibilities of perturbative methods and alternative approaches (like QCD sum rule methods or lattice calculations) are required; in practice, they appear in the so called hadronic distribution amplitudes (DA's) which describe, for each hadron (and independently of the particular process under consideration), how its momentum is shared among the valence constituents. Although a complete, formal proof of the validity of these factorization procedures is still lacking, at least for exclusive processes, there is enough theoretical work which supports these models [1–3].

As it should be clear, PQCD models acquire full validity only at very high momentum transfer; however, it is not yet clear what that means in practice (that is, at what Q^2 scale we expect these models to become reliable). This is a controversial point [4], even if very recently new developments seem to justify the applicability of the models also for not so high momenta [5]. On the other hand, we must not forget that all the cross sections for exclusive processes behave, at high Q^2, as an inverse power of Q^2. This power increases with the number of valence constituents involved in the process and makes it more difficult, from an experimental point of view, to distinguish and measure these increasingly rare events from the bulk of the inclusive processes. We may then summarize the situation as follows: on one hand, theoretical models are surely under better control at very high Q^2 but, on the other hand, almost all the experimental information presently at our disposal falls in a range of Q^2 which, while not completely out of reach of PQCD techniques, is not fully recognized as an ideal laboratory for perturbative models. In particular, it is not clear the role played by higher order corrections (which are reflected in several possible modifications of the basic PQCD models). Unfortunately, the implementation of these higher twist contributions is quite intricate. It follows from what we said that a complementary, theoretical, and phenomenological analysis of all the presently available experimental measurements would be of great help in clarifying and improving PQCD models. Thus it is very useful and important to undertake all the possible efforts in order to shed light on controversial points and improve our understanding of exclusive processes.

We should also bear in mind that presently only for a few, relatively simple processes, calculations have been performed. This is due to the increasing complexity of calculations when more and more hadrons (and, as a consequence, partonic constituents) are involved in the process. When compared with the experimental results (all of which, with the possible exception of the proton form factor, are at intermediate values of Q^2), these calculations show several successes but also some failures. Most of these failures can be attributed to violations of the so-called helicity selection rules [1] (which are a specific property of PQCD models, valid to all orders in the strong coupling constant perturbative expansion) and their overcoming requires the introduction of higher twist effects.

Several attempts have been made in order to implement the original PQCD models taking into account higher order contributions, like higher valence Fock states [6], transverse momentum effects [5], $L \neq 0$ angular momentum components in hadron wave functions [7], constituent quark mass effects [8–10], and diquark correla-

tions inside baryons [11].

For constituent quark mass effects, in particular, a number of calculations [8–10] have been performed for several experimentally observed processes, some of which are allowed in PQCD models while others are forbidden by the helicity selection rules. In these calculations the elementary hadron constituents are given a mass, $m_i = x_i m_H$, where m_H is the hadron mass and x_i is the (light-cone) fraction of the hadron momentum carried by the ith constituent, opportunely weighted (in the convolution integral) by the corresponding hadron distribution amplitude. This effective mass could take into account (in a global way) several higher order effects which have been neglected in the ordinary PQCD models.

Although from a formal point of view this approach requires further justifications, it offers a relatively simple, parameter free, means for practical calculations. The results obtained in this way can be compared with the lowest order PQCD results, and, as we shall also see in the following, lead to predictions for several effects that can be experimentally tested at the present time or in the near future (see also Ref. [10]).

A few years ago Carimalo [12] considered mass correction effects on the polar angular distribution of baryon-antibaryon pairs produced in the exclusive decay of the J/ψ. As it can easily be seen, PQCD models predict, due to the helicity selection rules, a distribution of the type $1 + \cos^2\theta_B$ [1], where θ_B is the polar angle which specifies the direction of motion of the produced baryon in the J/ψ rest frame. Although experimental results are available only for a few baryons [13,14] and are in some cases affected by large statistical errors, there are clear indications that the angular distributions behave rather as $1 + a_B \cos^2\theta_B$, with, e.g., $a_p = 0.62 \pm 0.11$, $a_\Lambda = 0.62 \pm 0.22$ [14]. Here a_B (≤ 1) is a factor which can be expressed from the helicity amplitudes for the decay process, as we shall see in detail in the next section.

As shown in Ref. [12], mass corrections can in principle explain why $a_B < 1$ and lead to a better agreement between theoretical predictions and experimental results. Furthermore, since the parameter a_B is given as a ratio of squared helicity amplitudes, it is independent of the exact value of the baryon decay constant, and of several "fine tuning" details of the models (for example, how to treat the strong coupling constant in the convolution integrals); as such, their effects may be neglected almost completely, as will be clarified by our explicit calculations.

The results of Ref. [12] are obtained using a nonrelativistic bound-state model both for the decaying J/ψ and the produced baryons. However, while this approximation is well grounded for heavy quark bound states, like the J/ψ, it can be questionable for light hadrons. In Ref. [12] the proton decay constant was consistently fixed in order to reproduce the $J/\psi \to p\bar{p}$ decay width. A different, widely used approach (see, e.g., Ref. [1,15] and references therein) is to tentatively fix once and for all the baryon decay constant from QCD sum rules and consider different models for the baryon distribution amplitudes, including the nonrelativistic, QCD-sum-rule-inspired and asymptotic ones. A lattice calculation of the baryon decay constant in the nucleon case is consistent, within the inherent systematic uncertainties of the models, with that of QCD sum rules [16]. In this context, the nonrelativistic DA seems to systematically underestimate absolute quantities (i.e., not obtained as ratios of amplitudes) like the decay widths for charmonium exclusive decays, or the hadron form factors, by two or even three orders of magnitude, when compared to available experimental results (see, e.g., Ref. [15]). It is then quite reasonable to expect that the use of more refined distribution amplitudes could lead to significant modifications in the predicted values for a_B.

There is also a more specific reason to believe that the dependence of a_B on the DA's may not be negligible: the expression of a_B is given as a function of squared helicity amplitudes with different values of the constituent helicities. Then, the use of distribution amplitudes, which, like those derived from QCD sum rules, seem to indicate an unusual sharing of the hadron momentum among its constituents, can substantially modify different helicity amplitudes.

Based on these motivations, in the rest of this paper we shall present a derivation of the parameter a_B, generalizing the results of Ref. [12] to the case of a generic hadron distribution amplitude. The only restriction on the DA is that it must satisfy general symmetry properties. This allows us to compare the results obtained using different DA's and possibly to get useful information on them, independently of the way the baryon decay constant has been fixed.

We wish also to recall that an alternative approach for the calculation of a_B has been proposed by Kada and Parisi. In Ref. [17] these authors evaluated the strong contribution to a_B for the octet baryons, in the framework of a quark-(scalar)diquark model for the baryon structure. Diquarks are another possible way of taking into account higher order corrections (in particular correlations between two valence quarks in the baryon). They have been applied, with good success, to several exclusive processes at intermediate values of Q^2 [11].

It was stressed by Carimalo and by Claudson et al. [18] and will also be argued in the following, that electromagnetic (em) corrections to a_B might be by no means negligible.

The main problem with the em corrections is that they involve the em form factors of the octet baryons, for which calculations including mass corrections have not yet been done. Only for the nucleon, in the nonrelativistic approximation, there is a theoretical evaluation [19]. In the other cases we are forced to give for these corrections estimates based on the available experimental information.

This paper is organized as follows. In Sec. II we derive a general expression for the angular distribution of the baryons produced in $J/\psi \to B\bar{B}$ decays. In Sec. III we concentrate on the strong contribution to a_B, which controls the $B\bar{B}$ angular distribution, giving a derivation of the helicity amplitudes required for its calculation. In Sec. IV we discuss in detail the results obtained and their dependence from some subtleties of the models that in general, while not modifying qualitatively the conclu-

sions, can substantially affect the numerical results, in particular for absolute quantities like decay widths. In Sec. V we give a detailed analysis of electromagnetic corrections, trying to estimate, when possible, upper bounds on the consequent overall modification of a_B both from experimental and theoretical information. Finally, our conclusions and future perspectives are discussed in Sec. VI.

II. DERIVATION OF THE BARYON ANGULAR DISTRIBUTION FOR $J/\psi \to B\bar{B}$ DECAYS

Let us consider a J/ψ particle, produced in e^-e^+ colliders with unpolarized beams, which subsequently decays into a baryon-antibaryon pair:

$$e^-e^+ \to J/\psi \to B\bar{B} \quad . \tag{1}$$

The spin density matrix of the J/ψ, in its rest frame [which is also the center of mass (cm) frame of colliding beams, with the electron moving along the positive z direction], has the following expression (see, e.g., Ref. [20]):

$$\rho_{MM'}(J/\psi) = \frac{1}{N} \sum_{\lambda_{e^-},\lambda_{e^+}} A^e_{M;\lambda_{e^-},\lambda_{e^+}} A^{e\,*}_{M';\lambda_{e^-},\lambda_{e^+}} \quad , \tag{2}$$

where the A^e's are the helicity amplitudes for the process $e^-e^+ \to J/\psi$, M is the z component of the J/ψ total angular momentum in its rest frame, and λ_{e^-}, λ_{e^+} are the helicities of the electron and the positron, respectively; N is a normalization factor, such that $\mathrm{Tr}[\,\rho\,] = 1$.

It is not difficult to show that

$$\rho_{MM'} = 0 \quad \text{if } M \neq M' \quad , \tag{3}$$

$$\rho_{11} = \rho_{-1,-1} = \frac{1}{2}\left(1 + 2\frac{m_e^2}{M_\psi^2}\right)^{-1/2} \cong \frac{1}{2} \quad , \tag{4}$$

$$\rho_{00} = 2\frac{m_e^2}{M_\psi^2}\left(1 + 2\frac{m_e^2}{M_\psi^2}\right)^{-1/2} \cong 0 \quad , \tag{5}$$

where m_e and M_ψ are the masses of the electron and the J/ψ, respectively. As for the second step in our process, the decay $J/\psi \to B\bar{B}$, we have the general relation [21]

$$d\Gamma(J/\psi \to B\bar{B}) = \frac{1}{8(2\pi)^5}\left(1 - 4\frac{m_B^2}{M_\psi^2}\right)^{1/2}$$
$$\times \sum_{M,\lambda_B,\lambda_{\bar{B}}} \rho_{MM} |A_{\lambda_B\lambda_{\bar{B}};M}|^2 d\Omega_B \quad , \tag{6}$$

where the A's are the helicity amplitudes for the decay of a J/ψ, with third component M of the total angular momentum $J = 1$, into a baryon-antibaryon pair with helicities λ_B and $\lambda_{\bar{B}}$, respectively; m_B is the mass of the produced baryons.

Due to the symmetry around the $\hat{z}$ axis, we can put $\varphi_B = 0$ in our calculations and integrating over φ_B we obtain

$$\frac{d\Gamma(J/\psi \to B\bar{B})}{d(\cos\theta_B)} = \frac{1}{8(2\pi)^4}\left(1 - 4\frac{m_B^2}{M_\psi^2}\right)^{1/2}$$
$$\times \sum_{M,\lambda_B,\lambda_{\bar{B}}} \rho_{MM} |A_{\lambda_B\lambda_{\bar{B}};M}|^2 \quad . \tag{7}$$

Apart from overall factors, independent of θ_B, this is the quantity measured by the MARKII and DM2 collaborations [13,14].

We know from first principles [20] that the amplitudes $A_{\lambda_B\lambda_{\bar{B}};M}$ have the following general structure:

$$A_{\lambda_B\lambda_{\bar{B}};M}(\theta_B,\varphi_B) = \tilde{A}_{\lambda_B\lambda_{\bar{B}}}\, d^1_{M,\lambda_B-\lambda_{\bar{B}}}(\theta_B)\exp(iM\varphi_B) \quad , \tag{8}$$

where the "reduced" amplitude $\tilde{A}_{\lambda_B\lambda_{\bar{B}}}$ is independent of M and the angular variables and the $d^{\,J}(\theta_B)$ are the usual rotation matrices. Using the parity properties [20] for the $A_{\lambda_B\lambda_{\bar{B}};M}$ (which imply that $\tilde{A}_{-+} = \tilde{A}_{+-}$, $\tilde{A}_{--} = \tilde{A}_{++}$) and for the spin density matrix $\rho(J/\psi)$ ($\rho_{-1,-1} = \rho_{1,1}$), Eq. (7) may be rewritten as follows:

$$\frac{d\Gamma(J/\psi \to B\bar{B})}{d(\cos\theta_B)} = \frac{1}{8(2\pi)^4}\left(1 - 4\frac{m_B^2}{M_\psi^2}\right)^{1/2}$$
$$\times \left\{|\tilde{A}_{+-}|^2(\rho_{11} + \rho_{00}) + 2|\tilde{A}_{++}|^2\rho_{11}\right\}$$
$$\times \left\{1 + a_B\cos^2\theta_B\right\} \quad , \tag{9}$$

where

$$a_B = \frac{\left\{|\tilde{A}_{+-}|^2 - 2|\tilde{A}_{++}|^2\right\}(\rho_{11} - \rho_{00})}{|\tilde{A}_{+-}|^2(\rho_{11} + \rho_{00}) + 2|\tilde{A}_{++}|^2\rho_{11}} \quad . \tag{10}$$

If, with good approximation [see Eqs. (4) and (5)], we take $\rho_{11} = 1/2$ and $\rho_{00} = 0$, we finally get the simplified expression

$$\frac{d\Gamma(J/\psi \to B\bar{B})}{d(\cos\theta_B)} = \frac{1}{16(2\pi)^4}\left(1 - 4\frac{m_B^2}{M_\psi^2}\right)^{1/2}$$
$$\times \left\{|\tilde{A}_{+-}|^2 + 2|\tilde{A}_{++}|^2\right\}$$
$$\times \left\{1 + a_B\cos^2\theta_B\right\} \quad , \tag{11}$$

where

$$a_B = \frac{|\tilde{A}_{+-}|^2 - 2|\tilde{A}_{++}|^2}{|\tilde{A}_{+-}|^2 + 2|\tilde{A}_{++}|^2} \quad . \tag{12}$$

We stress again that the results of Eqs. (11) and (12) only require the assumptions that the one-virtual-photon interaction dominates the J/ψ production process and that $m_e^2/M_\psi^2 \cong 0$.

Two further remarks are appropriate. (i) As it is clear from Eq. (8) and from the properties of the $d^{\,J}$,

$$\tilde{A}_{+-} = A_{+-;1}(\theta_B = \varphi_B = 0) \quad , \tag{13}$$

and

$$\tilde{A}_{++} = A_{++;0}(\theta_B = \varphi_B = 0) \quad . \tag{14}$$

Then it is sufficient to calculate the amplitudes $A_{+-;1}$, $A_{++;0}$ in the particular, convenient kinematic configuration $\theta_B = \varphi_B = 0$, in order to know, with the help of Eq. (8) and of the parity symmetry properties, all the amplitudes. (ii) As it will be explicitly shown in the next section, if $m_B = 0$ then $\tilde{A}_{++} = 0$ also; in this case we recover the old PQCD result of Brodsky and Lepage [22], that is, $a_B = 1$.

From Eq. (6) we easily get the expression of the total decay width for the process

$$\Gamma(J/\psi \to B\bar{B}) = \frac{1}{6(2\pi)^4}\left(1 - 4\frac{m_B^2}{M_\psi^2}\right)^{1/2}$$

$$\times \left\{|\tilde{A}_{+-}|^2 + |\tilde{A}_{++}|^2\right\} \quad . \tag{15}$$

III. THE STRONG CONTRIBUTION TO THE PARAMETER a_B: EVALUATION

In this section and in the next one we shall neglect the electromagnetic corrections to a_B, which will be considered in Sec. V, and concentrate on the strong contribution, which from now on will be called a_B^s. As we briefly sketched in Sec. I, the calculation of the helicity amplitudes for the physical process $J/\psi \to B\bar{B}$, $A_{\lambda_B \lambda_{\bar{B}};M}^s$ (the suffix s is a reminder that we are considering only the strong interaction in what follows), consists of several steps. First of all we need to calculate the amplitude for the hard interaction among the elementary (valence) constituents of the involved hadrons. To lowest order in the strong coupling constant, the only (topologically distinct) Feynman graph is shown in Fig. 1 (where the notation is also defined). All other possible graphs of the same
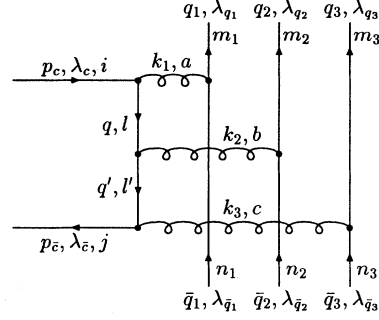

FIG. 1. The Feynman diagram which, to lowest order in α_s, describes the elementary process $Q\bar{Q} \to q_1 q_2 q_3 \bar{q}_1 \bar{q}_2 \bar{q}_3$, for a quarkonium state with charge conjugation $C = -1$. In the $Q\bar{Q}$ center-of-mass frame, $c^\mu = (E, \mathbf{k}/2)$ and $\bar{c}^\mu = (E, -\mathbf{k}/2)$, where $\mathbf{k}$ is the relative momentum between the c and $\bar{c}$ quarks; $q_i = x_i p_B$ and $\bar{q}_i = y_i p_{\bar{B}}$ ($i = 1, 2, 3$), with $p_B^\mu = (E, \mathbf{p}_B)$, $p_{\bar{B}}^\mu = (E, -\mathbf{p}_B)$, and $\mathbf{p}_B = (p\sin\theta_B\cos\varphi_B, p\sin\theta_B\sin\varphi_B, p\cos\theta_B)$. a, b, c, $i, j, l, l', m_{1,2,3}$, $n_{1,2,3}$ are color indices and the λ's label helicities.

order can be obtained from this one by a permutation of the final fermionic lines. We do not take into account explicitly all these graphs because their contribution is accounted for by opportunely choosing the final hadron wave functions. Without giving nonessential details of the intermediate steps of the calculation, we present directly the expression of this amplitude in the particular kinematic configuration $\theta_B = \varphi_B = 0$ (the relative momentum between the c and $\bar{c}$ quarks, $\mathbf{k}$, has also been set equal to zero. This procedure is proper when considering $L = 0$ bound states; for $L \neq 0$ the limit $\mathbf{k} \to \mathbf{0}$ should be taken in a subsequent step, i.e., after the integration over the angular part of the charmonium wave function).

$$T^s_{\lambda_{q_1}\lambda_{q_2}\lambda_{q_3},\lambda_{\bar{q}_1}\lambda_{\bar{q}_2}\lambda_{\bar{q}_3};\lambda_c\lambda_{\bar{c}}}(\mathbf{k} = \mathbf{0}; \theta_B = \varphi_B = 0)$$

$$= -16c_F g_s^6 \frac{1}{M_\psi^5} \frac{1}{\prod_{i=1}^3 \left[x_i y_i + (x_i - y_i)^2 \epsilon_B^2\right]}$$

$$\times \frac{1}{2x_1 y_1 - x_1 - y_1 + 2(x_1 - y_1)^2 \epsilon_B^2} \frac{1}{2x_3 y_3 - x_3 - y_3 + 2(x_3 - y_3)^2 \epsilon_B^2}$$

$$\times \left\{ \left[x_1 y_3 + x_3 y_1 + 2(x_1 - y_1)(x_3 - y_3)\epsilon_B^2\right] \left[\delta_{\lambda_{q_1},-\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},-\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},-\lambda_{\bar{q}_3}}\delta_{\lambda_{q_1},-\lambda_{q_2}}\delta_{\lambda_{q_2},-\lambda_{q_3}}\delta_{\lambda_{q_3},\lambda_c}\delta_{\lambda_c,-\lambda_{\bar{c}}} \right. \right.$$

$$+ \epsilon_B \left(\delta_{\lambda_{q_1},\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},-\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},-\lambda_{\bar{q}_3}}\delta_{\lambda_{q_3},\lambda_{q_2}}\delta_{\lambda_c,\lambda_{q_3}}\delta_{\lambda_c,-\lambda_c}\delta_{\lambda_c,\lambda_{\bar{c}}} + \delta_{\lambda_{q_1},-\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},-\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},\lambda_{\bar{q}_3}}\delta_{\lambda_{q_1},\lambda_c}\delta_{\lambda_{q_2},-\lambda_c}\delta_{\lambda_c,\lambda_{\bar{c}}}\right)$$

$$\left. + \epsilon_B^2 \delta_{\lambda_{q_1},\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},-\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},\lambda_{\bar{q}_3}}\delta_{\lambda_{q_2},\lambda_c}\delta_{\lambda_c,-\lambda_{\bar{c}}} \right] - \left[x_1 x_3 + y_1 y_3 - 2(x_1 - y_1)(x_3 - y_3)\epsilon_B^2\right]$$

$$\times \left[\epsilon_B \delta_{\lambda_{q_1},-\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},-\lambda_{\bar{q}_3}}\delta_{\lambda_{q_1},\lambda_c}\delta_{\lambda_{q_3},-\lambda_c}\delta_{\lambda_c,\lambda_{\bar{c}}} + \epsilon_B^2 \left(\delta_{\lambda_{q_1},\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},-\lambda_{\bar{q}_3}}\delta_{\lambda_{q_3},\lambda_c}\delta_{\lambda_c,-\lambda_{\bar{c}}} \right.\right.$$

$$\left. \left. \left. + \delta_{\lambda_{q_1},-\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},\lambda_{\bar{q}_3}}\delta_{\lambda_{q_1},\lambda_c}\delta_{\lambda_c,-\lambda_{\bar{c}}}\right) + \epsilon_B^3 \delta_{\lambda_{q_1},\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},\lambda_{\bar{q}_3}}\delta_{\lambda_c,\lambda_{\bar{c}}}\right]\right\} , \tag{16}$$

where c_F is the color factor which, once the convolution with the final hadron wave functions is made, takes the value $c_F = 5/(18\sqrt{3})$ and, as usual, $g_s = \sqrt{4\pi\alpha_s}$; $x_i(y_i)$ represents the (light-cone) fraction of the baryon (antibaryon) four-momentum carried by the ith quark (antiquark); $\epsilon_B = m_B/M_\psi$. It is clear from this equation the role played by mass corrections in allowing spin flips along the final fermionic lines. We see that higher ϵ_B powers correspond to terms where more spin flips are present. On the contrary, if we neglect quark masses Eq. (16) reduces, apart from a nonessential constant factor due to a different notation, to the result of Brodsky and Lepage [22]:

$$T^{(0)s}_{\lambda_{q_1}\lambda_{q_2}\lambda_{q_3},\lambda_{\bar{q}_1}\lambda_{\bar{q}_2}\lambda_{\bar{q}_3};\lambda_c\lambda_{\bar{c}}} = -16c_F g_s^6 \frac{1}{M_\psi^5} \frac{1}{\prod_{i=1}^{3} x_i y_i} \frac{1}{2x_1 y_1 - x_1 - y_1} \frac{1}{2x_3 y_3 - x_3 - y_3}$$

$$\times (x_1 y_3 + x_3 y_1)\delta_{\lambda_{q_1},-\lambda_{\bar{q}_1}}\delta_{\lambda_{q_2},-\lambda_{\bar{q}_2}}\delta_{\lambda_{q_3},-\lambda_{\bar{q}_3}}\delta_{\lambda_{\bar{q}_1},-\lambda_{q_2}}\delta_{\lambda_{\bar{q}_2},-\lambda_{q_3}}\delta_{\lambda_{\bar{q}_3},\lambda_c}\delta_{\lambda_c,-\lambda_{\bar{c}}} \quad . \tag{17}$$

We would like to stress that, in principle, the g_s^2 factors coming from different virtual gluons should be evaluated at the corresponding values of transferred Q^2 [1]. For the moment we do not discuss explicitly this problem, which will be analyzed in detail in the next section, where numerical results are presented.

Next we convolute the elementary helicity amplitude $T^s_{\{\lambda\}}$ (for brevity we indicate by $\{\lambda\}$ the collection of all the helicities from which an amplitude depends) with the final hadron wave functions:

$$M^s_{\lambda_B,\lambda_{\bar{B}};\lambda_c\lambda_{\bar{c}}} = \int [d\tilde{x}][d\tilde{y}]\psi_{B,\lambda_B}(\tilde{x})T^s_{\{\lambda\}}(\tilde{x},\tilde{y})\psi_{\bar{B},\lambda_{\bar{B}}}(\tilde{y}) , \tag{18}$$

where $\int [d\tilde{z}]$ stands for $\int_0^1 dz_1 dz_2 dz_3 \delta(1-z_1-z_2-z_3)$ and we used $\tilde{z}$ as a shorthand for (z_1, z_2, z_3). The amplitude $M^s_{\{\lambda\}}$ refers to the decay of a free c, $\bar{c}$ quark pair into the final baryon-antibaryon pair.

Once the amplitudes M^s have been evaluated (we will do this below), the final step consists in integrating these amplitudes over the proper $c\bar{c}$ bound state wave function, taken as usual in the nonrelativistic approximation, such that the physical amplitude $A^s_{\lambda_B,\lambda_{\bar{B}};M}$, to which we are interested in, is given by the following general expression:

$$A^s_{\lambda_B,\lambda_{\bar{B}};M} = \sum_{\lambda_c\lambda_{\bar{c}}} \left(\frac{2L+1}{4\pi}\right)^{1/2} C^{\frac{1}{2}\;\frac{1}{2}\;S}_{\lambda_c,-\lambda_{\bar{c}}\;\lambda} C^{L\;S\;J}_{0\;\lambda\;\lambda}$$

$$\times \int d^3k\, M^s_{\lambda_B,\lambda_{\bar{B}};\lambda_c\lambda_{\bar{c}}}(\mathbf{k})$$

$$\times D^{J\,*}_{M\lambda}(\beta,\alpha,0)\,\psi_C(\mathbf{k}) , \tag{19}$$

where the C's are the Clebsch-Gordan coefficients, $\lambda = \lambda_c - \lambda_{\bar{c}}$, $\mathbf{k} = (k,\alpha,\beta)$ is the relative momentum between the c and $\bar{c}$ quarks and finally $\psi_c(\mathbf{k})$ is the (momentum-space) charmonium wave function. In particular, for the J/ψ $L = 0$, so we can take from the beginning (in the full nonrelativistic approximation) $\mathbf{k} \rightarrow \mathbf{0}$ without loss of generality [this property has been used in the derivation of Eq. (16)]. Then we can see that

$$A^s_{+-;1} = \sqrt{2}\,\pi|R_s(0)|M^s_{+-;+-} , \tag{20}$$

$$A^s_{++;0} = 2\pi|R_s(0)|M^s_{++;++} , \tag{21}$$

where $R_s(0)$ is the value of the $L = 0$ charmonium wave function at the origin. By comparing the theoretical prediction for the decay width $\Gamma(J/\psi \rightarrow e^-e^+)$, $\Gamma_{ee} \cong (16/9)\alpha^2|R_s(0)|^2/M_\psi^2$, with the available experimental data [23] we can estimate $|R_s(0)| \simeq 0.737$ GeV$^{3/2}$. Here we have considered only the two amplitudes that, as it was discussed in the previous section, are sufficient to recover the expressions of all the others, when use is made of Eq. (8). In Eq. (21) the relation $M^s_{++;--} = M^s_{++;++}$ has also been used [the validity of this relation can be proved from Eqs. (16) and (18)].

As it is clear from Eq. (18), while the T^s amplitude is the same for all the baryon pairs considered, the M^s amplitudes are different for the different baryons, essentially because in general the spin-flavor component of their wave function changes. So, we cannot give a general, explicit expression for the amplitude M^s, but we must consider separately all the different cases. The total hadron wave function consists of a color part (which is the same for all the baryons and has been included in the definition of the color factor of the amplitude T^s), a spin-flavor component and a dynamical part which describes, in momentum space, how the baryon four-momentum is shared among its valence constituents: the distribution amplitude. In general the DA, which as we said previously is a nonperturbative quantity, allows for non-SU$_f$(3)-symmetric configurations, as has been shown by several studies performed with the help of QCD sum-rule techniques [15,24–28]. However, in the particular case of an SU$_f$(3)-symmetric DA, as we shall see, we obtain the same result for all the baryons and the only differences are the (experimental) values of the baryon masses [12].

In order to explicitly calculate the M amplitudes and their dependence from the DA's we need to take from the literature the available proposed models. QCD sum-rule results for octet baryon DA's exist at present only for the nucleon [15,24–28] and for the Σ^+, Ξ^-, and Λ [15]. Then, below we give explicitly the expression for the distribution amplitudes and for the required M^s amplitudes in these cases.

The most general wave functions for the baryons considered here are the following.

For p, n, Σ^+, Ξ^- baryons:

$$\psi_{B,\lambda_B}(\tilde{x}) = 2\lambda_B \frac{F_B}{4\sqrt{6}}\{\varphi_B(123)f_{1,\lambda_B}(1)f_{2,-\lambda_B}(2)f_{3,\lambda_B}(3) + \varphi_B(213)f_{1,-\lambda_B}(1)f_{2,\lambda_B}(2)f_{3,\lambda_B}(3)$$

$$-2R_B T_B(123)f_{1,\lambda_B}(1)f_{2,\lambda_B}(2)f_{3,-\lambda_B}(3) + (1 \longleftrightarrow 3) + (2 \longleftrightarrow 3)\} ; \tag{22}$$

for the Λ baryon:

$$\psi_{\Lambda,\lambda_\Lambda}(\tilde{x}) = 2\lambda_\Lambda \frac{F_\Lambda}{4\sqrt{6}}\{\varphi_\Lambda(123)u_{\lambda_\Lambda}(1)d_{-\lambda_\Lambda}(2)s_{\lambda_\Lambda}(3) - \varphi_\Lambda(213)u_{-\lambda_\Lambda}(1)d_{\lambda_\Lambda}(2)s_{\lambda_\Lambda}(3)$$

$$-2R_\Lambda T_\Lambda(123)u_{\lambda_\Lambda}(1)d_{\lambda_\Lambda}(2)s_{-\lambda_\Lambda}(3) + \text{all permutations of (1, 2, 3)}\} . \tag{23}$$

We have introduced the notation

$$R_B = \frac{F_B^T}{F_B} \quad , \tag{24}$$

where F_B and F_B^T are constants related to the value of the baryon wave function at the origin. In Eq. (22) $f_{1,2,3}$ are the flavors appropriate to the particular baryon considered ($f_{1,2,3} = uud$ for the proton, udd for the neutron, uus for the Σ^+, and ssd for the Ξ^-). Isospin symmetry properties impose several relations between $\varphi_B(\tilde{x})$ and $T_B(\tilde{x})$; in the case of the nucleon the relations $R_N = 1$, $2T_N(1,2,3) = \varphi_N(1,3,2) + \varphi_N(2,3,1)$ also hold (see Ref. [15] for further details). By insertion of these expressions and Eq. (16) in Eq. (18), we can, after some algebra, derive the following M^s amplitudes.

For N, Σ^+, Ξ^- baryons:

$$
\begin{aligned}
M_{++;++}^s = {} & \frac{F_B^2}{96}\epsilon_B \int [d\tilde{x}][d\tilde{y}]C_B\bigg(2D_B\Big\{\varphi_B(123)\varphi_B(213) - 2R_B[\varphi_B(321)T_B(321) \\
& + \varphi_B(312)T_B(132)]\Big\} + E_B\Big[\varphi_B(132)\varphi_B(312) - 4R_BT_B(123)\varphi_B(213)\Big] \\
& + 2E_B\epsilon_B^2\Big\{\varphi_B^2(123) + \varphi_B^2(213) + \varphi_B^2(312) + 2R_B^2[2T_B^2(123) + T_B^2(132)]\Big\}\bigg) ,
\end{aligned}\tag{25}
$$

$$
\begin{aligned}
M_{+-;+-}^s = {} & -\frac{F_B^2}{48} \int [d\tilde{x}][d\tilde{y}]C_B\bigg(D_B\Big[\varphi_B^2(123) + 2R_B^2T_B^2(132)\Big] \\
& + D_B\epsilon_B^2\Big[\varphi_B(132)\varphi_B(312) - 4R_BT_B(123)\varphi_B(213)\Big] + 2E_B\epsilon_B^2\Big\{\varphi_B(123)\varphi_B(213) \\
& - R_B\Big[2\varphi_B(123)T_B(123) + \varphi_B(312)T_B(132) + \varphi_B(132)T_B(132)\Big]\Big\}\bigg) .
\end{aligned}\tag{26}
$$

For the Λ baryon:

$$
\begin{aligned}
M_{++;++}^s = {} & \frac{F_\Lambda^2}{48}\epsilon_\Lambda \int [d\tilde{x}][d\tilde{y}]C_\Lambda\bigg(-2D_\Lambda\Big\{\varphi_\Lambda(123)\varphi_\Lambda(213) + 2R_\Lambda[\varphi_\Lambda(321)T_\Lambda(321) \\
& + \varphi_\Lambda(312)T_\Lambda(312)]\Big\} - E_\Lambda\Big[\varphi_\Lambda(132)\varphi_\Lambda(312) + 4R_\Lambda\varphi_\Lambda(213)T_\Lambda(213)\Big] \\
& + 2E_\Lambda\epsilon_\Lambda^2\Big\{\varphi_\Lambda^2(123) + \varphi_\Lambda^2(213) + \varphi_\Lambda^2(312) + 2R_\Lambda^2[2T_\Lambda^2(123) + T_\Lambda^2(312)]\Big\}\bigg) ,
\end{aligned}\tag{27}
$$

$$
\begin{aligned}
M_{+-;+-}^s = {} & -\frac{F_\Lambda^2}{24} \int [d\tilde{x}][d\tilde{y}]C_\Lambda\bigg(D_\Lambda\Big[\varphi_\Lambda^2(123) + 2R_\Lambda^2T_\Lambda^2(132)\Big] \\
& - D_\Lambda\epsilon_\Lambda^2\Big[\varphi_\Lambda(312)\varphi_\Lambda(132) - 4R_\Lambda\varphi_\Lambda(213)T_\Lambda(123)\Big] \\
& - 2E_\Lambda\epsilon_\Lambda^2\Big\{\varphi_\Lambda(123)\varphi_\Lambda(213) + 2R_\Lambda[\varphi_\Lambda(321)T_\Lambda(321) + \varphi_\Lambda(312)T_\Lambda(312)]\Big\}\bigg) .
\end{aligned}\tag{28}
$$

In Eqs. (25)–(28) the following concise notation has been used: in each $\varphi(i,j,k)\varphi(l,m,n)$, $\varphi(i,j,k)T(l,m,n)$, and $\varphi^2(i,j,k)=\varphi(i,j,k)\varphi(i,j,k)$ product the first term is a function of $\tilde{x}$, the second of $\tilde{y}$. Furthermore we have defined

$$D_B = x_1y_3 + x_3y_1 + 2(x_1 - y_1)(x_3 - y_3)\epsilon_B^2, \tag{29}$$

$$E_B = -(x_1x_3 + y_1y_3) + 2(x_1 - y_1)(x_3 - y_3)\epsilon_B^2, \tag{30}$$

$$
\begin{aligned}
C_B = {} & -\frac{2560\pi^3\alpha_s^3}{9\sqrt{3}M_\psi^5}\frac{1}{\prod_{i=1}^3[x_iy_i + (x_i - y_i)^2\epsilon_B^2]} \\
& \times \frac{1}{[2x_1y_1 - x_1 - y_1 + 2(x_1 - y_1)^2\epsilon_B^2]} \\
& \times \frac{1}{[2x_3y_3 - x_3 - y_3 + 2(x_3 - y_3)^2\epsilon_B^2]} .
\end{aligned}\tag{31}
$$

Let us finally stress that, as it was anticipated, all the $M_{++;++}^s$ amplitudes vanish if we take $m_B \to 0$ (this means that $a_B^s = 1$), as it must be, because in this case there is not in the model any mechanism which allows for spin flips in the quark-gluon vertices, thus forcing the final baryons to have opposite helicities.

IV. THE STRONG CONTRIBUTION TO THE PARAMETER a_B: NUMERICAL RESULTS

In the expressions presented in the previous section for the M^s amplitudes, from which a_B^s can be evaluated [see Eqs. (12), (16), (18), (20), and (21)] we are left with the unknown form of the distribution amplitudes, which contain all the hadronization dynamics of the partons of the dominant, valence light-cone Fock state. As we said previously, these DA's are highly nonperturbative in nature, so perturbative QCD cannot say about them much more than their general, formal solution (expressed as an

infinite sum of Appell polynomials with unknown, non-perturbative, coefficients) and their evolution in Q^2 [1]. This is by no means a trivial result, but it is not enough to give us reliable expressions for the DA at values of Q^2 presently accessible. For $Q^2 \to \infty$, only the first term in the formal expansion of the wave function survives, so we obtain a simple expression, the so called asymptotic distribution amplitude. The asymptotic DA again contains an unknown coefficient, which, however, may be extracted from the best experimental data at hand. The present status of nonperturbative methods (namely QCD sum rule and lattice calculations, in our specific case) allows us to calculate only a few low moments of the DA's. From these, a model expression for the DA is proposed by opportunely truncating the infinite, formal expansion of perturbative QCD and finding approximate values of the unknown expansion coefficients by fitting the calculated moments. There are, of course, several difficulties in this procedure. These are mainly reflected in an apparently strong dependence of the model DA from the number of available moments; in sizable differences among the DA's proposed by different groups; in discrepancies between the QCD sum rule and the lattice results. In our calculation, we will take a phenomenological approach, considering all the DA expressions available and analyzing how the results for a_B^s depend on each of them.

Let us first consider the case of the nucleon. There are several versions of QCD sum rule model DA's in the literature. For completeness we give these below, together with the nonrelativistic and the asymptotic DA's.

$$\varphi^{\mathrm{nr}}(\tilde{x}) = \prod_{i=1}^{3} \delta\left(x_i - \frac{1}{3}\right) , \qquad (32)$$

$$\varphi^{\mathrm{as}}(\tilde{x}) = 120 x_1 x_2 x_3 . \qquad (33)$$

The DA proposed by Chernyak and Zhitnitsky [24] is

$$\varphi_N^{\mathrm{CZ}}(\tilde{x}) = \varphi^{\mathrm{as}}(\tilde{x})[18.06 x_1^2 + 4.62 x_2^2$$
$$+ 8.82 x_3^2 - 1.68 x_3 - 2.94] . \qquad (34)$$

The DA of Chernyak, Ogloblin, and Zhitnitsky [15] is

$$\varphi_N^{\mathrm{COZ}}(\tilde{x}) = \varphi^{\mathrm{as}}(\tilde{x})[23.814 x_1^2 + 12.978 x_2^2$$
$$+ 6.174 x_3^2 + 5.88 x_3 - 7.098] . \qquad (35)$$

The DA of King and Sachrajda [26] is

$$\varphi_N^{\mathrm{KS}}(\tilde{x}) = \varphi^{\mathrm{as}}(\tilde{x})[20.16 x_1^2 + 15.12 x_2^2 + 22.68 x_3^2$$
$$- 6.72 x_3 + 1.68(x_1 - x_2) - 5.04] . \qquad (36)$$

That of Gari and Stefanis [27] is

$$\varphi_N^{\mathrm{GS}}(\tilde{x}) = \varphi^{\mathrm{as}}(\tilde{x})[-1.027 x_1^2 + 12.307 x_3^2 + 25.88 x_2$$
$$+ 111.32 x_1 x_3 + 9.105(x_1 - x_3) - 19.84] . \qquad (37)$$

The improved (heterotic) version of the Gari-Stefanis DA, proposed by Stefanis and Bergmann [28], is

$$\varphi_N^{\mathrm{het}}(\tilde{x}) = \varphi^{\mathrm{as}}(\tilde{x})[-19.773 + 32.756(x_1 - x_3)$$
$$+ 26.569 x_2 + 16.625 x_1 x_3$$
$$- 2.916 x_1^2 + 75.25 x_3^2] . \qquad (38)$$

In all these cases the nucleon decay constant F_N has roughly the same value:

$$|F_N| \cong 5.0 \times 10^{-3}\,\mathrm{GeV}^2 . \qquad (39)$$

Furthermore, as we said previously, $R_N = 1$ and $2 T_N(1, 2, 3) = \varphi_N(1, 3, 2) + \varphi_N(2, 3, 1)$ [see Eq. (22)].

The only parameter that must be fixed is the value of the strong coupling constant α_s. For the moment we will limit ourselves to take a fixed value of α_s for all three virtual gluons, using the perturbative expression for the running coupling constant:

$$\alpha_s(Q^2) = \frac{12\pi}{(11 n_c - 2 n_f) \ln(Q^2/\Lambda^2)} , \qquad (40)$$

where $n_c = 3$ is the number of colors and n_f is the number of active flavors (in this context, $n_f = 4$). From this point of view Q^2 must be set to an overall effective value, pertinent to the process considered, and we take the square of the J/ψ mass as this effective scale; we also use $\Lambda \cong 0.2$ GeV.

This is by no means the only possible way to treat α_s, as we shall see in the following. Of course, these ambiguities reflect some uncompletness of the theoretical models.

In order to evaluate the helicity amplitudes $A_{\lambda_B \lambda_{\bar{B}}; M}$ we finally need to compute quadruple integrals, which is done numerically. We have tested both our analytical results (the expressions of the amplitudes shown in the previous section) and our numerical calculations. For example, it is not difficult to see that if we consider a nonrelativistic DA we recover, through the different steps of our calculation, the analytical form of a_B^s proposed by Carimalo [12]. As for the numerical integrations, we have explicitly checked that if one of the usual representations of the Dirac function is inserted in the integrals, and a limit procedure is made for the parameter which enters the representation, the results smoothly tend to the analytical ones in the nonrelativistic case.

In Table I we give our results for a_N^s and for the total decay width $\Gamma^s(J/\psi \to N\bar{N}) = \Gamma^s_{N\bar{N}}$ for the nucleon.

TABLE I. The strong contribution to a_B and to the decay width $\Gamma(J/\psi \to B\bar{B})$ for the nucleon. Results obtained using the different nucleon DA's considered in the text are compared to experimental data.

DA	a_N^s	$10^7 \Gamma^s_{N\bar{N}}$ (GeV)
nr	0.688	0.002
as	0.667	0.026
CZ	0.561	0.587
COZ	0.565	0.826
KS	0.591	1.255
GS	0.963	0.168
het	0.689	1.671
MK2[a]	0.61 ± 0.23	1.85 ± 0.27
DM2[b]	0.62 ± 0.11	1.63 ± 0.38

[a]MARKII Collaboration, Ref. [13].
[b]DM2 Collaboration, Ref. [14].

Proton and neutron results differ in this case only for the small mass difference, so we do not present separate results (the situation will of course be different when we shall consider, in the next section, the electromagnetic corrections).

A comparison of our numerical results for the total decay widths with those of Refs. [12,15] needs some care. Reference [15] presents results only in the massless case and using $\alpha_s \simeq 0.3$ and $|R_s(0)| \simeq 0.690$ GeV$^{3/2}$. When opportunely rescaled to these parametrizations our massless results are in good agreement with those of Ref. [15]. Note, however, that from an analytical point of view our expression for $\Gamma(J/\psi \to B\bar{B})$ in the massless case differs from that of Ref. [15] by an overall factor of 80/81.

In the case of the nonrelativistic DA our analytical expression of $\Gamma(J/\psi \to B\bar{B})$ agrees with that of Ref. [12]. Let us stress, however, that our $F_N = 2^{5/2}F_N$ (Ref. [12]); that is, F_N (Ref. [12]) $\sim 5 \times 10^{-3}$ GeV^2 corresponds to our $F_N \sim 28 \times 10^{-3}$ GeV^2.

As can be seen from Table I, QCD sum rule DA's give results which are sizably different from those of the nonrelativistic case, even for an observable like a_B^s which is given as a ratio among squared helicity amplitudes. All the results (with the possible exception of the GS DA) are in agreement with the experimental measurements, but we must not forget that the experimental errors are large at present, and that electromagnetic contributions can in principle modify the theoretical values. We see also that there are non-negligible differences among the various QCD sum rule DA's. Of course the discrepancies are dramatically larger in the case of decay widths, so one can ask, why study the (relatively) little dependence of a_B^s on the DA? The point is that, as we shall see better in the following, the same reason which makes variations in a_B^s so small also tends to make a_B^s freer from several ambiguities of the model, giving, in our opinion, a more sensible test for the distribution amplitudes. However, in order to discriminate among different DA's we need better experimental results than those presently at our disposal. We hope they will be available in the near future. Note also that, with F_N given by Eq. (39), QCD sum rule DA's give results for the decay widths which are in rough agreement with the experiment, unlike the asymptotic and nonrelativistic ones (here, as in the following, no particular attempt in modifying parameters in

order to better reproduce experimental results has been made).

Apart from the nucleon, little is known about the form of the distribution amplitudes of the other octet baryons. To our knowledge, the only available models are those proposed by Chernyak et al. [15] for the Σ^+, Ξ^-, and Λ. From Ref. [15] we have for the Σ^+,

$$\varphi_\Sigma^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[0.36x_1^2 + 0.24x_2^2$$
$$+0.14x_3^2 - 0.54x_1x_2$$
$$-0.16x_3(x_1 + x_2) + 0.05(x_1 - x_2)] \quad , \quad (41)$$

$$T_\Sigma^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[0.32(x_1^2 + x_2^2) + 0.16x_3^2$$
$$-0.47x_1x_2 - 0.24x_3(x_1 + x_2)] \quad . \quad (42)$$

The decay constants are $|F_\Sigma| \simeq 5.1 \times 10^{-3}$ GeV^2, $|F_\Sigma^T| \simeq 4.9 \times 10^{-3}$ GeV^2.

For the Ξ^-,

$$\varphi_\Xi^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[0.38x_1^2 + 0.20x_2^2$$
$$+0.16x_3^2 - 0.26x_1x_2$$
$$-0.30x_3(x_1 + x_2) + 0.02(x_1 - x_2)] \quad , \quad (43)$$

$$T_\Xi^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[0.28(x_1^2 + x_2^2) + 0.18x_3^2$$
$$-0.16x_1x_2 - 0.35x_3(x_1 + x_2)] \quad , \quad (44)$$

with $|F_\Xi| \simeq 5.3 \times 10^{-3}$ GeV^2, and $|F_\Xi^T| \simeq 5.4 \times 10^{-3}$ GeV^2.

Finally for the Λ,

$$\varphi_\Lambda^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[0.44x_1^2 + 0.08x_2^2$$
$$+0.34x_3^2 - 0.56x_1x_2$$
$$-0.24x_3(x_1 + x_2) - 0.10(x_1 - x_2)] \quad , \quad (45)$$

$$T_\Lambda^{COZ}(\tilde{x}) = 42\varphi^{as}(\tilde{x})[1.2(x_2^2 - x_1^2) + 1.4(x_1 - x_2)] \quad (46)$$

and $|F_\Lambda| \simeq 6.3 \times 10^{-3}$ GeV^2, $|F_\Lambda^T| \simeq 6.3 \times 10^{-4}$ GeV^2.

In Table II, we compare the values of a_B^s obtained from these QCD sum rule DA's with those obtained in the nonrelativistic and the asymptotic case and with experimental results (note that, in the case of the Σ^+, experimental data are available for Σ^0 only; em corrections can in principle be different for Σ^+ and Σ^0, so a strict comparison of our results with the experimental data could be at this stage misleading). We see from Table II that

TABLE II. The strong contribution to a_B and to the decay width $\Gamma(J/\psi \to B\bar{B})$ for the Σ^+, Ξ^-, and Λ. Only the nonrelativistic, the asymptotic, and the COZ DA's are available for these particles.

| DA | Σ^{+a} | | Ξ^- | | Λ | |
	a_B^s	$10^7\Gamma_{B\bar{B}}^s$ (GeV)	a_B^s	$10^7\Gamma_{B\bar{B}}^s$ (GeV)	a_B^s	$10^7\Gamma_{B\bar{B}}^s$ (GeV)
nr	0.431	0.002	0.274	0.002	0.513	0.003
as	0.417	0.032	0.265	0.016	0.497	0.017
COZ	0.687	58.151	0.537	45.519	0.770	7.252
MK2[b]	0.7 ± 1.1	1.35 ± 0.35	-0.13 ± 0.55	0.97 ± 0.25	0.72 ± 0.36	1.35 ± 0.27
DM2[c]	0.22 ± 0.31	0.91 ± 0.27		0.60 ± 0.15	0.62 ± 0.22	1.18 ± 0.26

[a]Experimental results are available only for the $J/\psi \to \Sigma^0\bar{\Sigma}^0$ case.
[b]MARKII Collaboration, Ref. [13].
[c]DM2 Collaboration, Ref. [14].

for the Σ^+, Ξ^-, and Λ differences between nonrelativistic and QCD sum rule DA's results are also more marked. This is not unexpected, in that the presence of one or more valence s quarks breaks more severely the $SU_f(3)$ symmetry implicit in the nonrelativistic DA. Again, more precise measurements of a_B could be very useful in improving our understanding of (at least) the main features of the DA's. By the way, we stress that the decay widths for Σ^+, Ξ^-, Λ seem to be poorly reproduced, also when QCD sum rule DA's are used. Even if em corrections are not accounted for, it is unlikely that they can improve very much the results in this sense.

We want now to analyze better how our results depend on some ingredients of the models that, due to their (to some degree) effective nature, are not unambiguously fixed.

We basically concentrate on the behavior of the results when different ways of treating the strong coupling constant are considered. We stress that this is probably (together with the value of the baryon decay constant F_B) the main source of indeterminacy of the quantitative results of the model, once the DA has been fixed.

As we anticipated previously, when the calculations of the results quoted in Table I and Table II were discussed, the definition of α_s is in some sense ambiguous. We can make the following (reasonable but not derived from first principles) choices.

(i) Consider an overall, effective value of Q^2, Q_e^2, which sets the scale to which α_s must be evaluated, using Eq. (40), neglecting that in the model the momenta carried by the three virtual gluons in the hard scattering depend on $\tilde{x}, \tilde{y}$ and are different among them [1]. In particular, we choose $Q_e^2 \sim M_\psi^2$. Slightly different choices have been used in the literature (see, e.g., Refs. [15,24]); this is equivalent to an overall rescaling of the decay widths, while a_B is unchanged.

(ii) Take a running coupling constant depending on the squared momenta carried by the hard gluons, $\tilde{Q}_i^2 = Q_i^2(\tilde{x}, \tilde{y})$ (where i runs over the three gluons) [1]. However, if this choice is made, we face a serious problem: over the full field of variation of $\tilde{x}, \tilde{y}$, the $\tilde{Q}_i^2$ approach zero. There, then, it does not make sense to speak of perturbative calculations, since α_s increases more and more, invalidating any perturbative power expansion. Often

people prevent these problems by introducing an *ad hoc* cutoff, Q_0^2, such that

$$\alpha_s(Q^2) = \begin{cases} \dfrac{12\pi}{(11 n_c - 2 n_f) \ln(Q^2/\Lambda^2)} & \text{if } Q^2 > Q_0^2 \\ \alpha_0 = \alpha_s(Q_0^2) & \text{if } Q^2 \le Q_0^2 \end{cases},$$

(47)

where usually one takes $\alpha_0 = 0.3$ or 0.5. In some sense, this is a naïve way of accounting for nonperturbative effects which prevent α_s from taking higher and higher values [like Eq. (40) should imply], even if probably the exact behavior of α_s is only poorly mimed. Even if the $\tilde{Q}_i^2 \to 0$ region is suppressed by the behavior of the distribution amplitudes, this procedure is far from being satisfactory. Furthermore it introduces a dependence from the new parameter α_0.

(iii) Take into account in a more serious way nonperturbative effects, which may give rise to an effective mass for the gluons (see, e.g., Ref. [29]). This in turn modifies the gluon propagators and the perturbative expression for α_s, Eq. (40). To be rigorous, there are probably other minor modifications to be introduced into the fermion propagators and the quark-gluon vertices as well, if one consistently applies these methods. However, since we are only trying to estimate the dependence of our results from different ways of treating α_s, we limit ourselves to considering only the effective gluon mass effects on it. In this context, the expression of α_s to be used is of the type

$$\alpha_s(Q^2) = \dfrac{12\pi}{(11 n_c - 2 n_f) \ln\left[(Q^2 + 4 m_g^2)/\Lambda^2\right]},$$

(48)

where $m_g \sim 0.5$ GeV [29].

In Table III we show, for the nucleon case only (results for the other cases are very similar), how the three previous different choices of α_s influence our results for a_B^s and the decay widths. It is easy to see that a_B^s is very stable against these changes (as opposed to the total decay widths), giving us more confidence on the reliability of our estimates for a_B^s and on their usefulness.

Note that the nonrelativistic DA has not been taken into account, because in this case there is no ambiguity,

TABLE III. Dependence of the strong contribution to a_B and to $\Gamma(J/\psi \to B\bar{B})$ on different behaviors of the strong coupling constant α_s inside the convolution integral of Eq. (18), in the proton case. First column is the same as Table I, with α_s from Eq. (40) at $Q^2 = M_\psi^2$; in the 2nd and 3rd column use is made of Eq. (47), with $\alpha_0 = 0.3$ and 0.5, respectively; the 4th column presents the results when Eq. (48) is used.

| DA | $\alpha_s = 0.275$ | | $\alpha_s \le 0.3$ | | $\alpha_s \le 0.5$ | | $\alpha_s(m_g^2)$ | |
	a_N^s	$10^7 \Gamma_{N\bar{N}}^s$ (GeV)	a_N^s	$10^7 \Gamma_{N\bar{N}}^s$ (GeV)	a_N^s	$10^7 \Gamma_{N\bar{N}}^s$ (GeV)	a_N^s	$10^7 \Gamma_{N\bar{N}}^s$ (GeV)
as	0.667	0.026	0.667	0.044	0.666	0.466	0.667	0.210
CZ	0.561	0.587	0.561	1.144	0.564	10.013	0.561	5.716
COZ	0.565	0.826	0.565	1.369	0.567	12.078	0.564	6.894
KS	0.591	1.255	0.591	2.076	0.592	18.343	0.591	10.549
GS	0.963	0.168	0.963	0.275	0.954	1.917	0.963	1.160
het	0.689	1.671	0.686	2.763	0.682	24.245	0.691	14.020

in that the strong coupling constant is consistently fixed to the effective value $\alpha_s(M_\psi^2)$.

Let us conclude by stressing that we also must take into account the evolution of the DA with Q^2. It is usually said that this evolution has little effect on the results, in that its logarithmic behavior is masked by the stronger, powerlike behavior (in Q^2) of the elementary scattering amplitudes. Even if this may seem quite reasonable, we have explicitly checked for this assumption, particularly when modified running expressions for α_s have been adopted. As a matter of fact, while a_B is almost independent of the DA evolution, the modifications of the decay widths due to the evolution are not numerically negligible; however, they are not important from a qualitative point of view and do not change the results by more than a factor ~ 2, which is within the overall uncertainty we can expect, also in the more optimistic hypothesis, for absolute quantities.

V. ELECTROMAGNETIC CONTRIBUTIONS TO a_B

Up to now we have neglected the electromagnetic contributions to a_B. However, it has been shown that the corrections due to em processes can be sizable [12,18], and a careful analysis of a_B cannot avoid considering them. Unfortunately, as it was discussed by Carimalo [12], we are not able, at present, to give a full treatment of these effects. The reason is that the timelike em form factors of the baryons are involved, and there is no calculation of them including the mass corrections, as it has been done in this paper for the strong contribution. Only in the nucleon case, with a nonrelativistic DA, this calculation has been performed by Ji and Sill [19].

In our treatment we will follow, expand, and update the analysis made by Carimalo in Ref. [12]. After a general discussion of these corrections, we shall consider, from two different points of view, their estimation (strongly dependent on the baryon em form factors) and their contribution to a_B, limiting ourselves to the case of the nucleon. In fact, for the other octet baryons both experimental and theoretical knowledge of the form factors is too poor to allow a sensible phenomenological study.

We first try to give a full theoretical prediction for a_B in the framework of our model. As we said before, at present this is possible only for the nucleon and with a nonrelativistic distribution amplitude. However, we must not forget that nonrelativistic DA's give results far from the experimental ones for the em hadron form factors (which are generally underestimated and given, in some cases, with the wrong sign [1]). So, we must take the nonrelativistic case as the only one for which at present our model can give a full estimate of a_B, together with indications on the weight of em contributions, waiting for more reliable calculations of mass correction effects to baryon em form factors in the near future.

Next we shall consider these effects from a more phenomenological point of view, trying to estimate their order of magnitude from the available experimental results on the decay widths of the J/ψ into $B\bar{B}$ and e^-e^+ pairs,

respectively.

The reduced amplitudes defined in Eq. (8), including em contributions, may be written as follows:

$$\tilde{A}_{+\pm} = \tilde{A}_{+\pm}^s + \tilde{A}_{+\pm}^{em_1} + \tilde{A}_{+\pm}^{em_2} \quad , \tag{49}$$

where both em contributions are of order (α/α_s) with respect to the strong one, and all other higher order contributions in (α/α_s) are neglected. The term em_1 corresponds, for the hard scattering among the valence quarks, to Feynman graphs like those of Fig. 1, where a gluon is replaced by a photon; there are three topologically distinct graphs in this case, because we can replace each gluon by a photon. The contribution em_2 represents the decay of the J/ψ into a single virtual photon which in turn is directly (through the timelike em baryon form factors) coupled to the final baryon pair. Figure 2 shows the Feynman graphs pertinent to these terms.

Let us consider, first of all, the em_1 contribution. It is not difficult to see that, for the elementary, hard scattering amplitude, we have

$$T_i^{em_1} = -\frac{4}{5}\frac{\alpha}{\alpha_s}e_i T^s \quad , \tag{50}$$

where e_i is the electric charge, in units of the proton charge, of the ith quark, to which the virtual photon is coupled, and T^s is given by Eq. (16). As we said, there are three contributions like this, which sum up to give

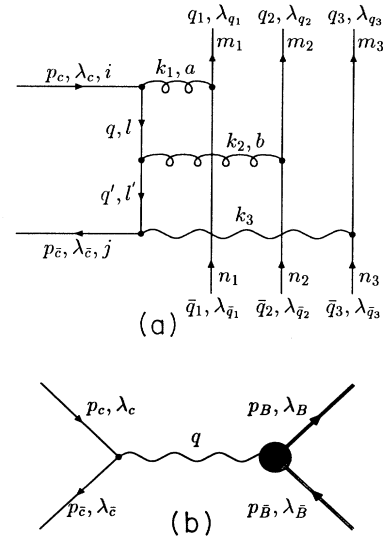$$T^{em_1} = -\frac{4}{5}\frac{\alpha}{\alpha_s}Q_B T^s = \delta_B T^s \quad . \tag{51}$$

FIG. 2. The two Feynman graphs which describe the leading (in α/α_s) electromagnetic corrections to the QCD lowest order term (see Fig. 1). (a) The contribution em_1, obtained by substituting one of the virtual gluons in Fig. 1 with a photon. There are two other contributions, obtained by replacing each of the other two gluons with a photon (see Fig. 1 for the notation). (b) The contribution em_2, coming from the direct coupling of a virtual photon produced in the $Q\bar{Q}$ decay with the final baryons.

Here $Q_B = e_1 + e_2 + e_3$ is the baryon electric charge (in units of the proton charge); we have also defined the constant δ_B ($\delta_B = 0$ for neutral baryons and $\delta_B \sim \pm 10^{-2}$ for $Q = \pm 1$ baryons). The subsequent steps which lead to the physical amplitudes $\tilde{A}$ are the same as for the strong contribution, so the same relation of Eq. (51) applies also to them:

$$\tilde{A}^{\text{em}_1}_{+\pm} = \delta_B \tilde{A}^s_{+\pm} \quad . \tag{52}$$

The case of the em$_2$ contribution is more subtle. First of all, given that the virtual photon couples directly to the final baryons, there is not an equivalent of the hard elementary amplitudes T, as in the preceding cases (or better, this step of the calculation is hidden in the baryon form factors). Then we must start from the decay of two free $c, \bar{c}$ quarks into the final baryon pair, which is described by means of the M amplitudes we defined, for the strong contribution, in Sec. III.

It can be seen that

$$M^{\text{em}_2}_{\lambda_B \lambda_{\bar{B}}; \lambda_c \lambda_{\bar{c}}}(\theta_B = \varphi_B = 0)$$

$$= \frac{8}{\sqrt{3}} \pi \alpha \frac{1}{M_\psi} \left\{ (F_1^B + k_B F_2^B) \right.$$
$$\times \left[2\delta_{\lambda_c \lambda_B} \delta_{\lambda_{\bar{c}}, -\lambda_{\bar{B}}} \delta_{\lambda_B, -\lambda_{\bar{B}}} + 2\epsilon_B \delta_{\lambda_c, \lambda_{\bar{c}}} \delta_{\lambda_B \lambda_{\bar{B}}} \right]$$
$$\left. + \frac{k_B}{2\epsilon_B} F_2^B (1 - 4\epsilon_B^2) \delta_{\lambda_c, \lambda_{\bar{c}}} \delta_{\lambda_B \lambda_{\bar{B}}} \right\} \quad , \tag{53}$$

where k_B is the baryon anomalous magnetic moment and $F_{1,2}^B$ are the well known Dirac and Pauli baryon form factors, respectively, at $q^2 = M_\psi^2$.

Then, from Eqs. (13), (14), (20), and (21) we obtain

$$\tilde{A}^{\text{em}_2}_{+-} = \frac{2^{9/2}}{\sqrt{3}} \pi^2 \alpha |R_s(0)| \frac{1}{M_\psi} G_M^B \quad , \tag{54}$$

$$\tilde{A}^{\text{em}_2}_{++} = \frac{2^5}{\sqrt{3}} \pi^2 \alpha |R_s(0)| \frac{1}{M_\psi} \epsilon_B G_E^B \quad , \tag{55}$$

where $G_M^B = F_1^B + k_B F_2^B$ and $G_E^B = F_1^B + (q^2/4m_B^2) k_B F_2^B$ are the Sachs baryon magnetic and electric form factors, respectively, again at $q^2 = M_\psi^2$.

We are now equipped with a formalism for calculating a_B including the leading electromagnetic corrections. Unfortunately a model consistent evaluation requires the analytical expressions for the baryon form factors including mass corrections, which at present are available only in the particular case of the nonrelativistic DA.

From the paper of Ji and Sill [19] we find

$$G_M^p(s) = 54 F_N^2 \pi^2 \alpha_s^2 \frac{1}{s^2} (-3 - 78\epsilon_p^2 + 136\epsilon_p^4) \quad , \tag{56}$$

$$G_E^p(s) = 54 F_N^2 \pi^2 \alpha_s^2 \frac{1}{s^2} (-11 - 24\epsilon_p^2 + 48\epsilon_p^4) \quad , \tag{57}$$

$$G_M^n(s) = 18 F_N^2 \pi^2 \alpha_s^2 \frac{1}{s^2} (9 + 140\epsilon_n^2 - 272\epsilon_n^4) \quad , \tag{58}$$

$$G_E^n(s) = 108 F_N^2 \pi^2 \alpha_s^2 \frac{1}{s^2} (7 - 10\epsilon_n^2) \quad . \tag{59}$$

It is easy to check that these results reproduce the corresponding massless results, when $\epsilon_{p,n} \to 0$ (see, e.g., Ref. [25] and references therein). However, we point out that mass corrections seem to be greater than expected, in that in some cases the higher order terms in $\epsilon_{p,n}^2$ are larger than the leading ones. We never encountered such a situation in our own calculations of mass correction effects.

By insertion of these results (at $s = M_\psi^2$) and Eq. (52) into Eq. (49) we can estimate the value of a_B and its variations with respect to the calculation made only with the strong contribution. The results are reported in Table IV, together with the values of the decay widths $\Gamma(J/\psi \to B\bar{B})$, both for the proton and the neutron. We see that percentual variations of a_B with respect to the pure strong contribution case are almost negligible: the effects of em contributions seem to be much smaller than those due to the choice of different DA's, see Table I.

Let us consider now what can be said on the em contributions from a phenomenological analysis, based upon the experimental information on some ingredients of the model.

From Eqs. (12), (49), (52), (54), and (55) we can see that

$$a_B = \frac{1 - \rho + a_B^s(1 + \rho)}{1 + \rho + a_B^s(1 - \rho)} \quad , \tag{60}$$

where we have defined

$$\rho = \left(\frac{1 + z x_{+-}}{1 + x_{+-}} \right)^2 \quad , \tag{61}$$

$$z = 2 \left(\frac{1 + a_B^s}{1 - a_B^s} \right)^{1/2} \epsilon_B \frac{G_E^B}{G_M^B} \quad , \tag{62}$$

and

$$x_{+-} = \frac{1}{1 + \delta_B} \frac{\tilde{A}^{\text{em}_2}_{+-}}{\tilde{A}^s_{+-}} \quad . \tag{63}$$

We see then that in order to estimate a_B we need to know a_B^s and the ratios G_E^B/G_M^B and x_{+-}. a_B^s has been evaluated in Secs. III and IV, while for the two ratios we must resort to experimental information.

Experimental estimates of x_{+-} can be obtained as follows: from Eqs. (15), (49), (54), and (55) we can write

TABLE IV. Theoretical predictions for a_B and the decay width $\Gamma_{B\bar{B}}(J/\psi \to B\bar{B})$, including em corrections, for the proton and the neutron. The 2nd column gives the percentual variation of a_B with respect to the pure strong contribution. Only the nonrelativistic DA is considered, and the results of Ji and Sill [19] for the nucleon em form factors are used.

	a_N	$\Delta a_N(\%)$	$10^7 \Gamma_{N\bar{N}}$(GeV)
p	0.696	1.2	0.002
n	0.677	-1.4	0.002

$$\Gamma_{B\bar{B}} = \Gamma_{ee}(1 - 4\epsilon_B^2)^{1/2}|G_M^B|^2 \left\{ \left(1 + \frac{1}{x_{+-}}\right)^2 \right.$$

$$\left. + 2\epsilon_B^2 \left|\frac{G_E^B}{G_M^B}\right|^2 \left(1 + \frac{1}{zx_{+-}}\right)^2 \right\} \quad , \quad (64)$$

where $\Gamma_{B\bar{B}} = \Gamma(J/\psi \to B\bar{B})$, $\Gamma_{ee} = \Gamma(J/\psi \to e^-e^+)$, and all other quantities have been previously defined. From Eq. (64) we can evaluate x_{+-}, once $|G_M^B|$ and the ratio G_E^B/G_M^B have been given. We take the experimental data on $\Gamma_{B\bar{B}}$, Γ_{ee} from Ref. [23].

Let us first consider the proton case. Very recently the timelike proton magnetic form factor at $q^2 \simeq 10$ GeV^2 has been measured [30], $|G_M^p(M_\psi^2)| = 0.026 \pm 0.002$ (experimental results on G_M^p were not available at the time Ref. [12] was published. By analytic continuation, it was assumed that to leading order in α_s, $G_M^p(q^2 = M_\psi^2) \simeq G_M^p(q^2 = -M_\psi^2) \simeq 0.012$. This does not seem to be the case [30]).

So, for the proton the only quantity we require which is not experimentally measured at present is the ratio $|G_E^B/G_M^B|$. One possibility is to assume approximate validity of the empirical relation $G_E^p/G_M^p \simeq 1/\mu_p$. However, there is nothing assuring that this relation can be verified in the kinematical regime of interest here. For this reason, we will subsequently take this ratio as a free parameter, allowing for reasonable variations of its value around the expected, empirical one, in order to study the dependence of a_B from this uncertainty.

Given that our results depend on the values of G_M^B and $|G_E^B/G_M^B|$ separately, let us clarify how measurements have been taken from Ref. [30]. It is seen from Ref. [30] that the measured value of $|G_M^p|$ depends on the value of the ratio $|G_E^p/G_M^p|$. This is so because there is not enough statistics in order to derive separate expressions of G_M^p and G_E^p. So, in our phenomenological analysis we proceed as follows: once the ratio G_E^p/G_M^p has been fixed to some value, we can estimate G_M^p from the experimental data [see Eq. (4) and Table I of Ref. [30]] and using the same fitting procedure of Ref. [30] to get the form factors at $s = M_\psi^2$. The first two columns of Table V give corrected values of a_p and percentual variation with respect to the pure strong contribution, when use is made of the

empirical relation $G_E^p/G_M^p \simeq 1/\mu_p$. We can see that em corrections can be relevant, and we probably cannot neglect them in a complete evaluation of a_B. We stress, however, that these estimates depend on the theoretical value of the strong contribution [see Eqs. (60)–(62)], as it can be seen from the spread of Δa_p over the various DA's considered. However, since we cover a large range of possibilities for the strong contribution, we think the order of magnitude of em corrections can be safely estimated to be of 5–10 %, for the proton. These estimates must be taken as upper limits.

Figure 3 shows the variation of a_p as a function of the ratio G_E^p/G_M^p and for some indicative DA's. The central value corresponds to the possible expected behavior, $G_E^p/G_M^p \simeq 1/\mu_p$.

We see from Fig. 3 that our previous conclusions are not greatly modified; only for big variations of the behavior of the G_E^p/G_M^p ratio from the empirical one, we can have relevant differences. Our conclusion is that the full range of variation of a_p with respect to a_p^s is of the order of 0–15 %.

In the case of the neutron the experimental information is also poorer than for the proton. We limit ourselves to extrapolate well known phenomenological behavior of the neutron form factors to the kinematic region $q^2 \simeq M_\psi^2$. From this point of view, we can at first assume $G_E^n \simeq 0$ and $G_M^n/\mu_n \simeq G_M^p/\mu_p$, at $q^2 = M_\psi^2$. The last two columns of Table V show the corresponding results for a_n and its percentual variation with respect to the pure strong contribution a_n^s. We see that this time em corrections induce a decrease of a_n, because G_M^n is negative. As for the absolute variation, this is slightly larger than in the proton case, and in the range of 9–15 %. Even if we do not report here the results, we have analyzed also
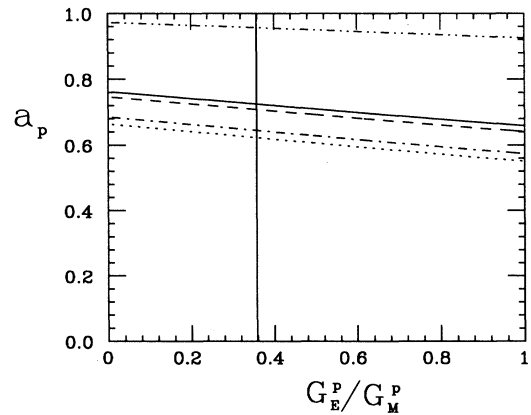
TABLE V. Results for $a_{p,n}$ obtained by adding to the theoretical values of $a_{p,n}^s$ phenomenological estimates of the em corrections. All the proposed DA's are considered. The $G_M^{p,n}(M_\psi^2)$ and $G_E^{p,n}(M_\psi^2)/G_M^{p,n}(M_\psi^2)$ values are fixed by using, when available, experimental data or by extrapolating them to the kinematical region under study (see text for details).

| DA | p | | n | |
---	a_p	$\Delta a_p(\%)$	a_n	$\Delta a_n(\%)$
nr	0.725	5.3	0.628	−8.9
as	0.708	5.9	0.605	−9.7
CZ	0.620	10.1	0.483	−14.7
COZ	0.623	9.9	0.487	−14.6
KS	0.645	8.8	0.517	−13.2
GS	0.956	−0.7	0.954	−0.8
het	0.726	5.3	0.629	−8.9

FIG. 3. The dependence of the full (QCD+QED) angular distribution parameter a_B [see Eq. (12)] from the poorly known G_E/G_M ratio, in the proton case. The vertical, solid line corresponds to the well known empirical behavior $G_E^p/G_M^p \simeq 1/\mu_p$. Different lines correspond to the various distribution amplitudes considered for the strong contribution (see text): nr (solid); as (dashed); COZ (dotted); KS (dot-dashed); GS (double dot-dashed). In this plot, the CZ and het DA results are almost indistinguishable from the COZ and nr ones, respectively.

for the neutron how em corrections change when G_M^n and G_E^n/G_M^n are allowed to vary around the empirical values given above. We find again a situation similar to that of the proton, possibly with a greater spread of variation.

VI. CONCLUSIONS

The study of exclusive processes involving hadrons at high energy scales, in the framework of perturbative QCD models, has made big improvements in recent years. However, several problems remain open and questionable, given that almost all the experimental information concerning exclusive processes is at intermediate energy scales. In such a situation, several higher order corrections may not be negligible and can contribute significantly, at least for those processes that are forbidden to lowest order in the PQCD models, but are experimentally well established. Unfortunately the implementation of these higher twist effects is very intricate, and only very recently significant progress has been made. Among these contributions, mass correction effects for the valence constituent quarks of the light hadrons involved are very promising. In fact, if from one hand a more rigorous theoretical explanation of their exact origin is required and auspicable, on the other hand their implementation is relatively easy and free of ambiguities: once their contribution is taken into account there are no free parameters to be fixed, and everything goes on without any further assumptions. A number of exclusive processes (in particular charmonium decays) have been analyzed in this framework in recent years, and several interesting consequences, experimentally testable at present or in the near future, have been proposed to check the validity of this model as compared to the PQCD ones or to alternative models including higher order effects. In this paper we have analyzed the effects of mass corrections on the angular distribution of octet baryon pairs created in J/ψ decays. Even if experimental measurements of these angular distributions have, for the moment, low statistics (at least in some cases), we hope they can be improved in the near future.

This was also the subject of an earlier paper by Carimalo. However, Carimalo used, as a first approach, a simplified, nonrelativistic model for the produced baryons. Over time it has been shown that the behavior of light hadrons, like the octet baryons, probably demands more accurate models of the baryon distribution amplitudes, like those we have considered here in a more general context. From a theoretical point of view, there are several reasons why the study of these angular distributions is interesting. First of all they are governed by a parameter which is given as a ratio of helicity amplitudes. As such, this parameter turns out to be quite independent of some details of QCD models presently at our disposal,

details which are not fully understood and can modify sizably the numerical results (albeit not the qualitative ones). Secondly, as we have shown, this parameter is sufficiently sensitive to the precise form of the distribution amplitudes to allow a discrimination among their main features, as soon as higher precision measurements will be available. Considering all the available models for the distribution amplitudes (mainly based on QCD sum rule calculations) we have shown indeed that the spread on the a_B values due to the change of the DA is of the order of 10–20 % in the nucleon case. Little is known about the DA's of the other octet baryons. Whenever possible, our calculations show an even larger dependence of a_B from the exact form of the DA, varying in the range of 20–50 %. Therefore a more accurate measurement of a_B could allow us to discriminate among different model DA's.

We have also considered in detail the role played by electromagnetic corrections. In fact, we cannot evaluate exactly these corrections at present. However, we give reasonable upper limit estimates of these contributions in the nucleon case, making use of all the possible available information, both theoretical and experimental.

Theoretical estimates (using a nonrelativistic approximation for the final baryon distribution amplitudes) suggest that em corrections could be negligible, both in the proton and neutron case.

On the contrary, we estimate from a phenomenological analysis that em corrections could be of the order of 5–10 % for the proton and of 10–15 % for the neutron. In the proton case em corrections tend to increase the value of a_B, while for the neutron they produce the opposite effect. In both cases, these corrections could be comparable to the modifications induced in the pure strong contribution a_B^s by different choices of the DA.

We can say very little about em corrections for the other octet baryons, since little is known about them, both experimentally and theoretically. So we limit ourselves to report, when possible, the pure strong contribution.

Of course, only more precise experimental information on the em baryon form factors could allow one to evaluate the exact contribution of em effects. Alternatively, the form factors could be evaluated in the framework of our model (including mass corrections), and a more consistent, theoretical result could be obtained. We leave this as a subject for future work.

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