Exclusive $J/\psi$ and $\psi'$ decays into baryon-antibaryon pairs

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Abstract
Within a perturbative approach we investigate decays of charmonium states into baryon-antibaryon pairs. Using a recently proposed wave function for the nucleon and suitable generalizations of it to the hyperons and decouplet baryons, we obtain the decay widths for the $B\bar{B}$ channels in reasonable agreement with data. An important difference to previous work is the use of the $c$-quark mass in the perturbative calculation instead of the charmonium mass. As a consequence of this feature our approach possesses, the $J/\psi$ and the $\psi'$ decay widths do not scale with a high power of the ratio of their masses.

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1 Introduction

The decay $J/\psi \rightarrow p\bar{p}$ has been investigated within perturbative QCD by Brodsky and Lepage first [1]. Later on this analysis has been repeated several times, e.g. [2, 3], and even extended to the $\Delta\Delta$ decay channel. It has been argued that the dominant dynamical mechanism is $c\bar{c}$ annihilation into three gluons and subsequent creation of light quark-antiquark pairs forming in turn the final state baryons. Three is the minimal number of gluons allowed in $J/\psi$ decays; $c\bar{c}$ annihilations through one or two gluons are forbidden by colour and $C$-parity, respectively. Contributions from annihilations through more than three gluons constitute higher order corrections. The dominance of annihilation through gluons is most strikingly reflected in the narrow width for hadronic channels in a mass region where strong decays have typically widths of hundreds of MeV [4]. The dynamical suppression at work here is perturbative QCD of higher orders (the total hadronic $J/\psi$ decay, for instance, is an $O(\alpha_s^3)$ process) and is customarily regarded as evidence for the Zweig rule. Since the $c$ and the $\bar{c}$ quarks only annihilate if they are separated by distances less than about $1/m_c$ and since the average virtuality of the gluons is about $1 \text{ GeV}^2$ one may expect perturbative QCD to be at work although corrections are presumably substantial. Indeed as the previous perturbative analyses, performed in the standard hard scattering approach, i.e. in collinear approximation, showed the $J/\psi$ decay into $p\bar{p}$ seems to be fairly well described. This is in marked contrast to the case of the nucleon form factor where soft physics seems to dominate in the experimentally accessible region of momentum transfer [5, 6, 7]. A point to criticize in these studies of the $J/\psi$ decays is the treatment of the strong coupling constant $\alpha_s$. Since, as we mentioned, the average virtuality of the gluons is about $1 \text{ GeV}^2$ one would expect $\alpha_s$ to be of the order of 0.4 to 0.5 rather than 0.2 to 0.3 as is customarily chosen [2, 3]. Since $\alpha_s$ enters to the sixth power into the expression for the width a variation of $\alpha_s$ from, say, 0.2 to 0.3 would lead to a change by a factor of 11 for the width. Thus, a large factor of uncertainty is hidden in these calculations preventing any severe test of the wave function utilized.

In constrast to previous works [1, 2, 3] we will not use the collinear approximation but rather the modified perturbative approach of Sterman et al. [8] in which transverse degrees of freedom are retained and Sudakov suppressions, comprising those gluonic radiative corrections not included in the evolution of the wave function, are taken into account. An important advantage of the modified perturbative approach is that the strong coupling constant can be used with a renormalization scale depending on the momentum fractions the quarks carry und thus large logs from higher orders of perturbative QCD are avoided. This choice of the renormalizaton scale entails singularities of $\alpha_s$ which are, however, compensated by the Sudakov factor. Hence, there is no uncertainty in the use of $\alpha_s$. This is to be contrasted with the standard perturbative approach where either $\alpha_s$ is evaluated at a renormalization scale that is a constant fraction of $M^2_\psi$, or at a momentum fraction dependent scale in which case $\alpha_s$ is to be “frozen” at a certain value (typically 0.5) in order to avoid uncompensated $\alpha_s$ singularities in the end-point regions. The modified perturbative approach possesses another interesting feature: the soft end-point regions are strongly suppressed. Therefore, the bulk of the perturbative contribution comes from regions where the internal quarks and gluons are far off-shell. In contrast to the nucleon form factor the $J/\psi \rightarrow B\bar{B}$ amplitude is not end-point sensitive. The suppression of the end-point regions does not, therefore, lead to a substantial reduction of the $J/\psi \rightarrow B\bar{B}$ amplitude. For the same reason, the size of that amplitude does not exhibit an extreme sensitivity to the baryon wave function utilized in the calculation as is, for instance, the
where \((2.1a)\) holds for all octet baryons except the \(\Lambda\) and the \(\Sigma^0\) for the decuplet baryons. Section 4 is devoted to the calculation of the few properties of light-cone wave functions. In section 3 we present the wave functions used in the analysis described in detail subsequently and briefly recapitulate a preliminary, experimental results for channels other than \(p\bar{p}\) have been reported [9].

The paper is organized as follows: In section 2 we introduce the octet baryon wave functions used in the analysis described in detail subsequently and briefly recapitulate a few properties of light-cone wave functions. In section 3 we present the wave functions for the decuplet baryons. Section 4 is devoted to the calculation of the \(J/\psi \to B_8\bar{B}_8\) decay widths. This analysis is extended to the decays into decuplet baryons (section 5) and to decays of other quarkonia into \(B\bar{B}\) pairs (section 6). Finally, section 7 contains our conclusions.

## 2 The wave functions of the octet baryons

Generalizing the ansatz made for the nucleon in [5, 6, 10], we write the valence Fock states of the lowest lying octet baryons as (the plane waves are omitted for convenience)

\[
|B_8, +\rangle = \frac{\varepsilon_{a_1a_2a_3}}{\sqrt{3!}} \int [dx][d^2k_\perp] \left\{ \Psi_{123}^{B_8} | f_{1+}^{a_1} f_{2+}^{a_2} f_{3+}^{a_3} \rangle + \Psi_{213}^{B_8} | f_{1+}^{a_1} f_{2+}^{a_2} f_{3+}^{a_3} \rangle \right. \\
- \left. \left( \Psi_{132}^{B_8} + \Psi_{231}^{B_8} \right) | f_{1+}^{a_1} f_{2+}^{a_2} f_{3+}^{a_3} \rangle \right\} \tag{2.1a}
\]

\[
|\Lambda, +\rangle = \frac{\varepsilon_{a_1a_2a_3}}{\sqrt{2}} \int [dx][d^2k_\perp] \left\{ \Psi_{123}^{\Lambda} | u_+^{a_1} d_+^{a_2} s_+^{a_3} \rangle - \Psi_{213}^{\Lambda} | u_+^{a_1} d_+^{a_2} s_+^{a_3} \rangle \\
+ \left( \Psi_{132}^{\Lambda} + \Psi_{231}^{\Lambda} \right) | u_+^{a_1} d_+^{a_2} s_+^{a_3} \rangle \right\} \tag{2.1b}
\]

where \((2.1a)\) holds for all octet baryons except the \(\Lambda\) and the \(\Sigma^0\). Obviously, for the proton and the \(\Sigma^+\) \(f_1\) represents an \(u\) quark and \(f_2\) either a \(d\) or a \(s\) quark, respectively. For the \(\Xi^-\) \(f_1\) represents a \(s\) quark and \(f_2\) a \(d\) one. The states of the neutron, \(\Sigma^-\) and \(\Xi^-\) are obtained from those of the proton, \(\Sigma^+\) and \(\Xi^-\) states, by exchanging \(u \leftrightarrow d\), respectively. The baryon is assumed to be moving rapidly in the 3-direction. Hence the ratio of transverse, \(k_{\perp i}\), to longitudinal momenta, \(x_i p_i\), of the quarks is small and one may still use a spinor basis on the light cone. The integration measures are defined by

\[
[dx] \equiv \prod_{i=1}^3 dx_i \delta(1 - \sum_i x_i) \quad [d^2k_\perp] \equiv \frac{1}{(16\pi)^2} \prod_{i=1}^3 d^2k_{\perp i} \delta^{(2)}(\sum_i k_{\perp i}) \tag{2.2}
\]
The quark $f_i$ is characterized by the momentum fraction $x_i$, by the transverse momentum $k_{\perp i}$ as well as by its helicity $\lambda_i$ and colour $a_i$. A three-quark state is then given by
\[
|f_{1\lambda_1}^a f_{2\lambda_2}^a f_{3\lambda_3}^a\rangle = \frac{1}{\sqrt{x_1 x_2 x_3}} |f_1^{a_1}; x_1, k_{\perp 1}, \lambda_1\rangle |f_2^{a_2}; x_2, k_{\perp 2}, \lambda_2\rangle |f_3^{a_3}; x_3, k_{\perp 3}, \lambda_3\rangle.
\] 
(2.3)

The single-quark states are normalized as
\[
\langle f_i^{a_i'}; x_i', k_{\perp i}', \lambda_i' | f_i^{a_i}; x_i, k_{\perp i}, \lambda_i\rangle = 2x_i(2\pi)^3 \delta_{a_i'a_i} \delta_{f_i f_i} \delta_{\lambda_i' \lambda_i} \delta(x_i' - x_i) \delta^{(3)}(k_{\perp i}' - k_{\perp i}).
\] 
(2.4)

Since the 3-component of the orbital angular momentum, $L_3$, is assumed to be zero the quark helicities sum up to the baryon’s helicity. (2.1a) is the most general ansatz for the $L_3 = 0$ projection of the three-quark nucleon wave function [11]. From the permutation symmetry between the two $u$ quarks and from the requirement that the three quarks have to be coupled in an isospin 1/2 state it follows that there is only one independent scalar wave function. If the $L_3 \neq 0$ projections are included the entire nucleon state is described by three independent functions [11]. In general there are more than one scalar wave function for the other octet baryons if SU(3)$_F$-symmetry breaking is taken into account.

As already expressed in (2.1) we nevertheless assume that each octet baryon is described by a single scalar wave function which, for convenience, we write as
\[
\Psi_{123}^{B_8}(x, k_{\perp}) = \frac{1}{8\sqrt{3!}} f_{(8)}(\mu_F) \phi_{123}^{B_8}(x, \mu_F) \Omega_{(8)}(x, k_{\perp}).
\] 
(2.5)

We assume that SU(3)$_F$ symmetry is only broken by a quark mass dependence of $\phi_{B_8}$ (see below). $f_{(8)}$, being related to the wave function at the origin of the configuration space, is identified with the nucleon parameter $f_N$ whose value was determined in [6] to amount to $6.64 \cdot 10^{-3}$ GeV$^2$ at the scale of reference $\mu_0 = 1$ GeV.

The transverse momentum dependence of the baryon wave function is parameterized by a simple symmetric Gaussian
\[
\Omega_{(8)}(x, k_{\perp}) = (16\pi^2)^2 \frac{a_{(8)}^2}{x_1 x_2 x_3} \exp \left[ -a_{(8)}^2 \sum_{i=1}^3 k_{\perp i}^2/x_i \right]
\] 
(2.6)

and the transverse size parameter $a_{(8)}$ is assumed to be the same for all octet baryons. A value of 0.75 GeV$^{-1}$ is used for that parameter (see [6]).

The last item to be specified in (2.5) is the distribution amplitude, $\phi_{ijk}^{B_8}(x, \mu_F) \equiv \phi_{B_8}(x_i, x_j, x_k, \mu_F)$, of an octet baryon $B_8$, which is conventionally normalized to unity
\[
\int [dx] \phi_{123}^{B_8}(x, \mu_F) = 1.
\] 
(2.7)

The distribution amplitude representing the wave function integrated over transverse momenta up to the factorization scale, $\mu_F$, can be expanded upon the eigenfunctions of the evolution kernel being linear combinations of Appell polynomials (see [12, 13])
\[
\phi_{123}^{B_8}(x, \mu_F) = \phi_{AS}(x) \left[ 1 + \sum_{n=1}^{\infty} B_n^{B_8}(\mu_F) \tilde{\phi}_{123}^n(x) \right]
\] 
(2.8)

where $\phi_{AS}(x) \equiv 120 x_1 x_2 x_3$ is the asymptotic distribution amplitude [12]. Evolution is incorporated by the factorization scale dependences of $f_{(8)}$ and the expansion coefficients $B_n$:
\[
f_{(8)}(\mu_F) = f_{(8)}(\mu_0) \left( \frac{\ln(\mu_0/\Lambda_{QCD})}{\ln(\mu_F/\Lambda_{QCD})} \right)^{2/3\beta_0}, \quad B_n^{B_8}(\mu_F) = B_n^{B_8}(\mu_0) \left( \frac{\ln(\mu_0/\Lambda_{QCD})}{\ln(\mu_F/\Lambda_{QCD})} \right)^{\gamma_n/\beta_0}
\] 
(2.9)
\begin{table}
\begin{tabular}{|l|l|l|}
\hline
\(n\) & \(\tilde{\phi}^n_{123}(x)\) & \(\tilde{\gamma}_n\) \\
\hline
1 & \(x_1 - x_3\) & \(20/9\) \\
2 & \(-2 + 3(x_1 + x_3)\) & \(8/3\) \\
3 & \(2 - 7(x_1 + x_3) + 8(x_1^2 + x_3^2) + 4x_1x_3\) & \(32/9\) \\
4 & \((x_1 - x_3)(1 - 4/3(x_1 + x_3))\) & \(40/9\) \\
5 & \(2 - 7(x_1 + x_3) + 14/3(x_1^2 + x_3^2) + 14x_1x_3\) & \(14/3\) \\
\hline
\end{tabular}
\end{table}

Table 1: Eigenfunctions and reduced anomalous dimensions for helicity 1/2 baryons.

where \(\beta_0 \equiv 11 - 2/3 n_f\). The exponents \(\tilde{\gamma}_n\) are the reduced anomalous dimensions. Because they are positive fractional numbers increasing with \(n\) [12], higher order terms in (2.8) are gradually suppressed. The reduced anomalous dimensions and the eigenfunctions \(\tilde{\phi}^n_{123}\) are listed in Tab. 1 where the notation of [14] is adopted.

In [6] the nucleon distribution amplitude was found to have the simple form

\[
\phi^N_{123}(x, \mu_0) = \phi_{AS}(x) \left[ 1 + \frac{3}{4} \tilde{\phi}^1_{123}(x) + \frac{1}{4} \tilde{\phi}^2_{123}(x) \right] = 60 x_1 x_2 x_3 \left[ 1 + 3 x_1 \right].
\]

The nucleon wave function is fully specified now and, before turning to the discussion of the hyperon distribution amplitudes, we note in passing that the probability of the nucleon’s valence Fock state is 0.17.

A suitable hyperon distribution amplitude is constructed by taking (2.10) and, in order to incorporate the empirically known breaking of SU(3)\(_F\) symmetry, multiplying it with a factor

\[
\exp \left( -a^{(8)}_m m_s^2 \frac{x_j}{x_j} \right)
\]

whenever the quark \(j\) is a strange one. That factor bears resemblance to the BHL exponential [15]. \(a^{(8)}\) is the transverse size parameter already introduced in (2.6) and \(m_s\) is a still to be adjusted parameter related to the strange quark mass. Explicitly our hyperon distribution amplitudes read

\[
\phi^\Sigma_{123}(x, \mu_0) = N_{\Sigma} \phi^N_{123}(x) \exp \left( -a^{(8)}_m m_s^2 \frac{x_3}{x_3} \right),
\]

\[
\phi^\Lambda_{123}(x, \mu_0) = \frac{1}{3} N_{\Lambda} \left( \phi^N_{123}(x) + 2 \phi^{N}_{321}(x) \right) \exp \left( -a^{(8)}_m m_s^2 \frac{x_3}{x_3} \right)
\]

\[
\phi^\Xi_{123}(x, \mu_0) = N_{\Xi} \phi^N_{123}(x) \exp \left( -a^{(8)}_m m_s^2 \left[ \frac{1}{x_1} + \frac{1}{x_2} \right] \right).
\]

The constants \(N_{B_s}\) ensure the correct normalizations (see (2.7)) of the hyperon distribution amplitudes \((N_{B_s} = 1\) for \(m_s = 0\)). The particular combination of \(\phi^N\)’s appearing in the \(\Lambda\) case (2.12b) is required by SU(3)\(_F\) symmetry.

In order to take into account evolution properly we expand the distribution amplitudes (2.12a)-(2.12c) upon the eigenfunctions of the evolution kernel (see (2.8)) up to terms of
<table>
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<th></th>
<th>Set 1 ((m_s = 0))</th>
<th>Set 3 ((m_s = 350 \text{ MeV}))</th>
<th></th>
<th>Set 2 ((m_s = 150 \text{ MeV}))</th>
<th>Set 4 ((m_s = 480 \text{ MeV}))</th>
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<tr>
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<td></td>
<td>0.000</td>
<td>-0.282</td>
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<td>0.721</td>
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<tr>
<td>(\Xi)</td>
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<td>0.000</td>
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</table>

Table 2: Expansion coefficients \(B_n(\mu_0)\) of octet baryon distribution amplitudes for various values of the parameter \(m_s\).

order \(n = 5\). In Tab. 2 we quote four sets of expansion coefficients \(B_n\) corresponding to the following scenarios: Set 1 is obtained from \(m_s = 0\), set 2 from \(m_s = 150 \text{ MeV}\) (current strange quark mass) and set 4 from \(m_s = 480 \text{ MeV}\) (constituent strange quark mass). The intermediate set 3 corresponds to \(m_s = 350 \text{ MeV} \approx (480^2 - 330^2)^{1/2} \text{ MeV}\) (difference between the squares of strange and light constituent quark masses). This value appears if (2.11) is interpreted as the ratio of the BHL exponentials for a hyperon and the nucleon distribution amplitude.

In Fig. 1 we show contour plots of the four octet baryon distribution amplitudes for \(m_s = 480 \text{ MeV}\) in order to illustrate the effect of the mass exponential (2.11). It can be seen that, compared to the nucleon case, the maxima of the \(\Sigma\) and \(\Lambda\) distribution amplitudes are shifted to the right, i.e. to larger \(x_3\) values, whereas that of the \(\Xi\) distribution amplitude is shifted to the left. On the average, \(s\) quarks carry larger momentum fractions than \(d\) quarks if \(m_s > 0\).

Distribution amplitudes for the octet baryons are also discussed in [16]. As compared to our ones these QCD sum rule based distribution amplitudes are strongly concentrated in the end-point regions and exhibit a more significant breaking of \(SU(3)_F\) symmetry. Moreover, they reveal three pronounced maxima near the end-points. In summary, the distribution amplitudes proposed in [16] are very different from ours and also from the asymptotic distribution amplitude at experimentally accessible scales. As our distribution amplitudes but to a much greater extent, they possess the property that, on the average, a \(u\) \((s)\) quark in the proton or the \(\Sigma^+\) \((\Xi^-, \Lambda)\) with the same helicity as its parent baryon carries the largest fraction of the baryon momentum.
3 The wave functions of the decuplet baryons

In analogy to (2.1) we write the valence Fock states of the decuplet baryons $\Delta$ and $\Sigma^*$ with helicity $\lambda_{B_{10}} = 1/2$ as

$$
|\Delta^{++}, + \rangle = \frac{\varepsilon_{a1a2a3}}{\sqrt{2}} \int [dx][d^2k_\perp] \Psi_{123}^{\Delta} | u_+^{a_1} u_-^{a_2} u_+^{a_3} \rangle
$$

$$
| B_{10}, + \rangle = \frac{\varepsilon_{a1a2a3}}{\sqrt{3!}} \int [dx][d^2k_\perp] \left\{ \Psi_{123}^{B_{10}} | f_1^{a_1} f_2^{a_2} f_3^{a_3} \rangle + \Psi_{213}^{B_{10}} | f_1^{a_1} f_2^{a_2} f_3^{a_3} \rangle + \Psi_{132}^{B_{10}} | f_1^{a_1} f_2^{a_2} f_3^{a_3} \rangle \right\}.
$$

For the $\Delta^+$ ($\Sigma^{*-}$) $f_1$ represents an $u$ ($d$) quark and $f_2$ a $d$ ($s$) one. Again, the states of the $\Delta^0$, $\Delta^-$ and $\Sigma^{**}$ are obtained from those of the $\Delta^+$, $\Delta^{++}$ and $\Sigma^{*-}$ by exchanging $u \leftrightarrow d$, respectively. The generalization to the other helicity states is trivial. There is obviously only one independent scalar wave function in the case of the $\Delta$ and we assume that the same shall apply to the $\Sigma^*$. The scalar wave functions $\Psi_{123}^{B_{10}}$ of the decuplet baryons $B_{10}$ ($= \Delta, \Sigma^*$) are parameterized in a fashion similar to the octet baryon case:

$$
\Psi_{123}^{B_{10}}(x, k_\perp) = f_{(10)}(\mu_F) \phi_{123}^{B_{10}}(x, \mu_F) \Omega_{(10)}(x, k_\perp).
$$

$f_{(10)}$ is assumed to be equal for all members of the baryon decuplet. We again adopt the form (2.6) for the transverse momentum dependent part of the wave function, $\Omega_{(10)}$, with the transverse size parameter $a_{(8)}$ replaced by $a_{(10)}$.

---

1 We omit the discussion of the decuplet states $\Xi(1530)$ and $\Omega^-$ since the perturbative approach is not applicable to the decays $J/\psi \rightarrow \Xi(1530)\Xi(1530)$ due to almost zero momentum transfer. The decay $J/\psi \rightarrow \Omega^- \bar{\Xi}^+$ is even kinematically forbidden. The decay into $\Sigma^*$ is the borderline case for the application of the perturbative approach.
Decuplet baryons in helicity $1/2$ states have the same eigenfunctions of the evolution kernel and the same anomalous dimensions as the octet baryons (see Tab. 1). Evolution of the helicity $3/2$ baryons is, on the other hand, different [12]. We refrain from giving details on that case here since we do not consider such baryons.

In the particular case of the $\Delta$ it is tempting to use a completely permutation symmetric distribution amplitude since the $\Delta$ is composed of are three light quarks in symmetric spin and flavour states. Noting that the permutation symmetric part of the nucleon distribution amplitude (2.10) is just equal to the asymptotic distribution amplitude, we take $\phi^{\Delta}_{123}(x) = \phi_{\text{AS}}(x)$ and construct $\phi^{\Sigma^*}(x)$ analogously to $\phi^{\Sigma}(x)$

$$
\phi_{123}^{\Sigma^*}(x) = N_{\Sigma^*} \phi_{\text{AS}}(x) \exp \left( -\frac{a_{(10)}^2 m_N^2}{x_3} \right) .
$$

The evolution of the $\Sigma^*$ distribution amplitude is treated as in the case of the octet baryons by expanding (3.3) upon the eigenfunctions $\tilde{\phi}_{123}^n(x)$ up to $n = 5$. In distinction from the octet baryon case, the QCD sum rule based distribution amplitudes for the decuplet baryons [13, 17] are not very different from the ones we are proposing.

### 4 Decays of the $J/\psi$ into octet baryons

Now, with the model wave functions at hand, we can calculate the decay width of $J/\psi$ into octet baryon-antibaryon pairs within the modified perturbative approach. The helicity amplitudes of these processes may be decomposed covariantly to read

$$
\mathcal{M}^{B_8}_{\lambda_1 \lambda_2 \lambda} = \pi_{B_8}(p_1, \lambda_1) \left[ \mathcal{B}^{B_8} \gamma_{\mu} + \mathcal{C}^{B_8} \frac{(p_1 - p_2)_{\mu}}{2 m_{B_8}} \right] v_{B_8}(p_2, \lambda_2) \epsilon^\mu(\lambda)
$$

where $m_{B_8}$ is the mass of an octet baryon. $p_1$ ($p_2$) and $\lambda_1$ ($\lambda_2$) denote the momentum and the helicity of an octet baryon (antibaryon), respectively. $u_{B_8}$ and $v_{B_8}$ are their spinors (normalized as $\pi_{B_8} u_{B_8} = 2 m_{B_8}$) and $\epsilon$ is the polarization vector of the $J/\psi$. In a leading twist perturbative approach the helicity amplitudes are determined by the invariant $\mathcal{B}^{B_8}$ ($\mathcal{C}^{B_8} = 0$). Implicitly this ensures hadronic helicity conservation. The $J/\psi \rightarrow B_8 \bar{B}_8$ decay width reads

$$
\Gamma(J/\psi \rightarrow O\bar{O}) = \frac{\rho_{\text{p.s.}}(m_{B_8}/M_{J/\psi})}{48 \pi M_{J/\psi}} \sum |\mathcal{M}^{B_8}_{\lambda_1 \lambda_2 \lambda}|^2
$$

in the $J/\psi$ rest frame ($M_{J/\psi}$ is the $J/\psi$ mass). If the $J/\psi$ is produced in $e^+e^-$ annihilations it is transversely polarized with respect to the beam direction and the angular distribution of the baryons emitted in the $J/\psi$ rest frame then exhibits a $1 + \cos^2 \theta$ dependence (up to corrections of $O(m_{B_8}^2/M_{J/\psi}^2)$) characteristic of perturbative QCD [1]. The function $\rho_{\text{p.s.}}$ in (4.2) is the usual phase space factor

$$
\rho_{\text{p.s.}}(z) = \sqrt{1 - 4z^2}.
$$

As we said in the introduction we are going to calculate the invariant $\mathcal{B}^{B_8}$ within the modified perturbative approach proposed in [8], thus generalizing our analysis of the $J/\psi$ decay into nucleon-antinucleon pairs [6]. As in previous perturbative calculations [1, 2, 3] the $J/\psi$ meson is treated as a non-relativistic $c\bar{c}$ system and $O(\nu^2/c^2)$ corrections are
neglected. In contrast to the decays of $P$-wave charmonium \[18\] the $c\bar{c}g$ Fock state is suppressed in exclusive decays by inverse powers of the large scale which is provided by the $c$-quark mass $m_c$ \[19, 20\], relative to the $c\bar{c}$ Fock state and is, therefore, neglected in our analysis. The use of $m_c$, strictly speaking $2m_c$, as the large scale rather than the charmonium mass is consistent with the neglect of relativistic corrections. It is also well in the spirit of a perturbative approach since in the internal $c$-quark propagators the $c$-quark mass appears. The $J/\psi$ state is written in a covariant fashion

$$|J/\psi; q, \lambda\rangle = \frac{\delta_{ab}}{\sqrt{3}} \left( \frac{f_\psi}{2\sqrt{6}} \right) \frac{1}{\sqrt{2}} (\hat{q} + M_\psi) f(\lambda) \quad (4.4)$$

where $a$ and $b$ are colour indices and $f_\psi$ is the $J/\psi$ decay constant being related to the $J/\psi$ wave function at the origin of the configuration space. Merely that part of the wave function is, to a reasonable approximation, required in the calculation of $\mathcal{B}$ since, as we already mentioned in the introduction, the $c$ and the $\bar{c}$ quark only annihilate if their mutual distance is less than about $1/m_c$ \[4\] which is smaller than the $J/\psi$ radius \[21\]. $M_\psi$ is replaced by $2m_c$ in the calculation of $\mathcal{B}_\psi$ and the baryon masses are neglected. Only the phase space factor in (4.2) is evaluated with the physical masses, $M_\psi$ and $m_{B_8}$.

The invariant $\mathcal{B}_8^{B_8}$ receives its dominant contribution from the graphs with three intermediate gluons, see Fig. 2. Within the modified perturbative approach the three-gluon contribution $\mathcal{B}_{3g}^{B_8}$ to the $J/\psi$ decay into $B_8\bar{B}_8$ is of the form

$$\mathcal{B}_{3g}^{B_8} = \frac{f_\psi}{2\sqrt{6}} \int [dx][dx'] \int \frac{d^2b_1}{(4\pi)^2} \frac{d^2b_3}{(4\pi)^2} \hat{T}_H(x, x', b) \exp[-S(x, x', b, 2m_c)] \quad (4.5)$$

$$\times \left[ \hat{\Psi}_{123}(x, b)\hat{\Psi}_{123}(x', b) + \frac{1}{2} (\hat{\Psi}_{123}(x, b) + \hat{\Psi}_{321}(x, b)) (\hat{\Psi}_{123}(x', b) + \hat{\Psi}_{321}(x', b)) \right]$$

for $B_8 = N, \Sigma, \Xi$. This contribution is the same for all octet baryons belonging to the same isospin multiplet. Since the form of the $\Lambda$ Fock state (2.1b) is slightly different the three-gluon contribution reads

$$\mathcal{B}_{3g}^{\Lambda} = \sqrt{\frac{3}{2}} \frac{f_\psi}{2} \int [dx][dx'] \int \frac{d^2b_1}{(4\pi)^2} \frac{d^2b_3}{(4\pi)^2} \hat{T}_H(x, x', b) \exp[-S(x, x', b, 2m_c)] \quad (4.6)$$

$$\times \left[ \hat{\Psi}_{123}(x, b)\hat{\Psi}_{123}(x', b) + \frac{1}{2} (\hat{\Psi}_{123}(x, b) - \hat{\Psi}_{321}(x, b)) (\hat{\Psi}_{123}(x', b) - \hat{\Psi}_{321}(x', b)) \right]$$

in the case of the $\Lambda$. The representation of $\mathcal{B}_{3g}$ as a convolution of wave functions and a hard scattering amplitude $\hat{T}_H$ can formally be derived by using the methods described in detail by Botts and Sterman \[8\]. The $b_i$, canonically conjugated to the transverse momenta $\mathbf{k}_{\perp i}$, are the quark separations in the transverse configuration space. $\mathbf{b}_1$ and $\mathbf{b}_3$ correspond to the locations of quarks 1 and 3 in the transverse plane relative to quark 2 and $\mathbf{b}_2 = \mathbf{b}_1 - \mathbf{b}_3$. $\hat{\Psi}_{ijk}$ represents the Fourier transform of the wave function $\Psi_{ijk}$ (see Sect. 2).

$\hat{T}_H$ is the Fourier transform of the usual momentum space hard scattering amplitude to be calculated from the Feynman graphs shown in Fig. 2a. Up to corrections of order
Figure 2: Feynman graphs for the $J/\psi$ decay into baryon-antibaryon. a) Three-gluon contribution (graphs with permutated gluon lines are not shown), b) electromagnetic contribution.

$$\alpha_s^4, m_{B_b}^2/(4m_c^2)$$ and $b^2/(4m_c^2)$ the hard scattering amplitude in b space reads

$$\hat{T}_H(x,x',b) = -\frac{5}{27} 2^{12} f_\psi m_c^6 \frac{(x_1x'_2 + x_3x'_1)}{(\tilde{q}_i^2 + \tilde{g}_i^2)(\tilde{q}_3^2 + \tilde{g}_3^2)} \left( \prod_{i=1}^3 \alpha_s(t_i) \right) \int d^2b_0$$

$$\times \left[ \frac{i\pi}{2} H_0^{(1)}(\tilde{q}_1|b_1 + b_0|) - K_0(\tilde{q}_1|b_1 + b_0|) \right] \frac{i\pi}{2} H_0^{(1)}(\tilde{g}_2b_0)$$

$$\times \left[ \frac{i\pi}{2} H_0^{(1)}(\tilde{g}_3|b_3 + b_0|) - K_0(\tilde{q}_3|b_3 + b_0|) \right].$$

(4.7)

The quantities

$$\tilde{q}_i^2 = 2[x_i (1 - x'_i) + (1 - x_i) x'_i] m_c^2, \quad \tilde{g}_i^2 = 4x_i x'_i m_c^2$$

(4.8)

represent the virtualities of the internal quarks and gluons at zero transverse momenta. Since, as shown in [6], the hard scattering amplitude only depends on the sum of transverse momenta, $k_{\perp i} + k'_{\perp i}$, the transverse separation of any two quarks inside the baryon is the same as that of the corresponding antiquarks inside the antibaryon ($b_i = b'_i$). Physically, this property of the hard scattering amplitude means that baryon and antibaryon are created with identical transverse configurations of the quarks and antiquarks, respectively.

The auxiliary variable $b_0$ in (4.7) serves as a Lagrange multiplier to the constraint $\sum_i k_{\perp i} + k'_{\perp i} = 0$. Since the virtualities of the gluons are timelike, $\hat{T}_H$ includes complex-valued Hankel functions $H_0^{(1)}$ that are related to the usual modified Bessel functions $K_0$, appearing for space-like propagators, by analytic continuation.

The Sudakov factor $\exp[-S]$ entering (4.5) and (4.6) takes into account those gluonic radiative corrections not accounted for in the QCD evolution of the wave function as well as the renormalization group transformation from the factorization scale $\mu_F$ to the renormalization scales $t_i$ at which the hard amplitude $\hat{T}_H$ is evaluated. The Sudakov factor, originally derived by Botts and Sterman [8] and later on slightly improved, can be found for instance in [22]. The renormalization scales $t_i$ are defined in analogy to the case of electromagnetic form factors [8] as the maximum scale of either the longitudinal momentum or the inverse transverse separation associated with each of the gluons

$$t_1 = \max(\tilde{q}_1, \tilde{g}_1, 1/b_3), \quad t_2 = \max(\tilde{g}_2, 1/b_2), \quad t_3 = \max(\tilde{g}_3, \tilde{g}_3, 1/b_1).$$

(4.9)

Infrared cut-off parameters $\tilde{b}_i$ appear in the Sudakov factor which are naturally related to, but not uniquely determined by the mutual separations of the three quarks [23].
Figure 3: Percental accumulation of the three-gluon contribution to the width of the $NN$ channel from regions of internal momenta where $\prod_{i=1}^{3} \alpha_s(t_i) < \alpha_s^{\text{crit}} ^{3}$.

Following [5] we chose $\tilde{b}_i = \tilde{b} = \max \{b_1, b_2, b_3\}$. With this “MAX” prescription the three-gluon contribution $B_{3g}$ is unencumbered by $\alpha_S$ singularities in the soft end-point regions. As a consequence of the regularizing power of the “MAX” prescription, the perturbative contribution saturates in the sense that the results become insensitive to the inclusion of the soft regions. A saturation as strong as possible is a prerequisite for the self-consistency of the perturbative approach. The infrared cut-off $\tilde{b}$ marks the interface between the non-perturbative soft gluons, which are implicitly accounted for in the baryon wave function, and the contributions from soft gluons, incorporated in a perturbative way in the Sudakov factor. Obviously, the gliding factorization scale to be used in the evolution of the wave function, has to be chosen as $\mu_F = 1/\tilde{b}$.

As an inspection of (4.5), (4.6) and (4.7) reveals, a nine dimensional numerical integration has to be performed.\textsuperscript{2} Although this is a rather involved technical task it can be carried through with sufficient accuracy if some care is put into it. The numerical results are obtained from the wave functions discussed in Sect. 2 and for the following values of the $J/\psi$ decay constant and the $c$-quark mass: $f_{\psi} = 409$ MeV, $m_c = 1.5$ GeV. We evaluate $\alpha_s$ in the one-loop approximation with $n_f = 4$ and $\Lambda_{QCD} = 210$ MeV [24].

Before turning to the detailed discussion of the numerical results for the decay widths, we want to focus on an important feature of the modified perturbative approach. As mentioned in the introduction, one of the motivations for including the transverse hadronic structure and the Sudakov factor in the analysis is to achieve a theoretically self-consistent calculation in the sense that the bulk of the perturbative contribution is accumulated in regions where the strong coupling constant $\alpha_s$ is sufficiently small. A method to check whether or not this is the case is to set the integrand in (4.5) or (4.6) equal to zero in those regions where $\prod_{i=1}^{3} \alpha_s(t_i) > (\alpha_s^{\text{crit}})^3$ and to evaluate the three-gluon contribution to the decay width as a function of $\alpha_s^{\text{crit}}$. In Fig. 3 we show the accumulation profile for the nucleon case; it is typical of all baryons. Consulting Fig. 3, one sees that almost the entire result is accumulated in the comparatively narrow region of $\alpha_s$ between 0.4 and 0.6. The regions with $\prod_{i=1}^{3} \alpha_s(t_i) < 0.469^3$ provide 50 % of the total result. Hence, our calculation of the $J/\psi$ decay widths into octet baryon-antibaryon pairs is theoretically self-consistent.

\textsuperscript{2}Taking into account relativistic corrections to the $J/\psi$ wave function, i.e. its transverse momentum dependence, one would have to perform a 14 dimensional numerical integration which seems impossible with present day computers to a sufficient degree of accuracy.
Table 3: Results for the decay widths of $J/\psi$ into octet baryon-antibaryon pairs for the four sets of distribution amplitudes defined by the expansion coefficients $B_n$ quoted in Tab. 2 ($f_\psi = 409$ MeV, $m_c = 1.5$ GeV). For comparison we also quote the experimental results and, in the column labelled $\phi_{AS}$, predictions evaluated with distribution amplitudes constructed from the asymptotic distribution amplitude instead from (2.10) (with $m_s = 350$ MeV).

<table>
<thead>
<tr>
<th>Channel</th>
<th>Set 1</th>
<th>Set 2</th>
<th>Set 3</th>
<th>Set 4</th>
<th>$\phi_{AS}$</th>
<th>$\Gamma_{exp}$ [eV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$p\bar{p}$</td>
<td>174</td>
<td>174</td>
<td>174</td>
<td>174</td>
<td>140</td>
<td>186 ± 14</td>
</tr>
<tr>
<td>$\Sigma^0\Sigma^0$</td>
<td>136</td>
<td>128</td>
<td>113</td>
<td>108</td>
<td>97.8</td>
<td>110 ± 15</td>
</tr>
<tr>
<td>$\Lambda\bar{\Lambda}$</td>
<td>140</td>
<td>133</td>
<td>117</td>
<td>107</td>
<td>99.7</td>
<td>117 ± 14</td>
</tr>
<tr>
<td>$\Xi^-\Xi^+$</td>
<td>107</td>
<td>92.8</td>
<td>62.5</td>
<td>47.4</td>
<td>46.9</td>
<td>78 ± 18</td>
</tr>
</tbody>
</table>

In Tab. 3 we present our results for the $J/\psi \rightarrow B_8\bar{B}_8$ decay widths using the distribution amplitudes discussed in Sect. 2. For the sake of comparison we also expose the available experimental data [24]³. As can be seen from the results obtained with the distribution amplitudes termed set 1 the phase space factor $\rho_{p.s.}$ is an important but not sufficient element for the suppression of the hyperon channels. Since $m_s = 0$ for these distribution amplitudes the differences in the predictions for the widths, except for the $\Lambda\bar{\Lambda}$ case, are only due to $\rho_{p.s.}$. For $m_s > 0$ the additional suppression of the end-point regions leads to smaller decay widths for the hyperon channels. In order to demonstrate the strength of that reduction we display the invariant $B_{3g}^{B_8}$ versus $m_s$ in Fig. 4. The number of strange quarks embodied in a given baryon is reflected in differently strong $m_s$ dependences of $B_{3g}^{B_8}$. As inspection of Tab. 3 brings to view, the phase space corrected three-gluon contributions nicely reproduce the experimental pattern of the decay widths provided they are computed with the distribution amplitudes of set 3. The value of 350 MeV for the parameter $m_s$ used in the construction of these distribution amplitudes appears reasonable, considering the interpretation of the mass factor (2.11) as the BHL exponential [15]. Note that, as a consequence of the use of the $c$-quark mass in the calculation of $B_{3g}^{B_8}$, the result for the decay width of the $p\bar{p}$ channel given in Tab. 3 differs from that one reported in [6].

The distribution amplitudes defined by (2.10-2.12c) exhibit a little asymmetry with respect to permutations of the $x_i$. In order to study the importance of this asymmetry we also show in Tab. 3 results computed with the asymptotic distribution amplitude for the nucleon and with modifications of it for the hyperons that are constructed analogously to the distribution amplitudes described in Sect. 2. The asymmetry in the distribution amplitudes is seen to increase the magnitudes of the decay widths by about 30 % for the $\Xi^-\Xi^+$ channel and about 20% for the other channels while the pattern of the predictions remains unchanged.

Let us now assess the uncertainties of our predictions. The value of the $J/\psi$ decay constant used by us is determined from the leptonic $J/\psi$ decay width. Since the $J/\psi$

³We use PDG averages throughout. The original data are from [25, 26, 27, 28].
decay constant, or more generally, the decay constant, $f_n$, of a $n^3S_1$ quarkonium state is defined by

$$\langle 0 | j^\mu_{em} | n^3S_1 \rangle = f_n M_n \epsilon^\mu,$$  \hspace{1cm} (4.10)

the leptonic decay width of a $n^3S_1$ state reads

$$\Gamma(nS \rightarrow e^+e^-) = \frac{4\pi e_Q^2 \alpha^2 f_n^2}{3 M_n}$$  \hspace{1cm} (4.11)

where $e_Q$ is the charge of the heavy quark the charmonium state consists of. Reexpressing (4.11) in terms of the non-relativistic wave function at the origin of the configuration space, one arrives at the famous van Royen-Weisskopf width [29]. The use of the decay constant instead of the wave function at the origin implies that we relate the $J/\psi \rightarrow BB$ (or the $n^3S_1 \rightarrow BB$) widths to the leptonic width. By that means the uncertainties in the determination of the wave function at the origin via the usual van Royen-Weisskopf width cancel to a large extend. These uncertainties arise from relativistic and QCD corrections which seem to be large [30] but are not well known [21, 31].

The next uncertainty to be mentioned arises from the choice of the $c$-quark mass value. In accordance with calculations of the charmonium spectrum within non-relativistic potential approaches [21] and with a global fit of charmonium parameters [32] we take 1.5 GeV as the favoured value. That value has, for instance, also been used in a recent analysis of $P$-wave charmonium decays into two pions [18]. In spite of this, little changes of the $m_c$ value cannot be excluded and lead to an approximate rescaling of the decay widths by the factor $(1.5 \text{ GeV}/m_c)^8$.

The value of $\Lambda_{QCD}$ is also subject to uncertainties. A change of that value by, say, $\pm 20$ MeV which roughly represents the inaccuracy of our present knowledge of $\Lambda_{QCD}$ [24], would alter the theoretical decay widths by about $\pm 25\%$. We stress that for any changes of the $m_c$ and $\Lambda_{QCD}$ values the ratios of any two decay widths calculated by us remain approximately unchanged. This assertion does not only refer to the $J/\psi \rightarrow B_s \overline{B}_s$ decay widths but it also applies to the still to be discussed $J/\psi \rightarrow B_{10} \overline{B}_{10}$ and $\psi' \rightarrow BB$ widths.

In addition to the three-gluon contribution $B_{3g}^{B_8}$ there is a subdominant, although in some cases perhaps sizeable, isospin-violating electromagnetic one, $B_{em}^{B_8}$ [33], arising from
the graph shown in Fig. 2b. This contribution is proportional to the time-like magnetic form factor of the baryon. For the proton $B^p_{em}$ amounts to about 15% (in absolute value) of the total $B^p$ as is estimated from the recent measurement of the proton magnetic form factor in the time-like region [34]. Since the relative phase between $B^p_{3g}$ and $B^p_{em}$ is unknown we cannot simply add the two contributions. Therefore, we merely can state that, depending on the value of the relative phase, the electromagnetic contribution (including the interference term between it and $B^p_{3g}$) to the $J/\psi \rightarrow p\bar{p}$ decay width can be as large as 30% or only 2%. The comparison of our result with data indicates a relative phase close to $\pm \pi/2$ in the $p\bar{p}$ case.

At this point a remark concerning the $n\bar{n}$ decay channel is in order. Since, as said repeatedly, the three-gluon contribution respects isospin symmetry any difference between the $p\bar{p}$ and $n\bar{n}$ decay widths must be due to the electromagnetic contribution. From experiment it is known that the widths for $J/\psi \rightarrow p\bar{p}$ and $J/\psi \rightarrow n\bar{n}$ decays agree within the experimental errors [28] and that the time-like form factors of the proton and the neutron are approximately equal in modulus at least at $s = 5.4$ GeV$^2$ [35]. Thus, one may conclude that the relative phases between the three-gluon and the electromagnetic contributions are the same (up to a possible sign) for the proton and the neutron channel.

The size of the electromagnetic contribution to the hyperon channels may be estimated from a recent analysis of the octet baryon form factors within a diquark model [36]. With the help of a few rather well determined parameters that model is able to describe a large number of exclusive observables. In particular relevant for the present work is the prediction that, in the space-like region, the magnetic form factors of the $\Sigma^+$ and $\Sigma^-$ have opposite signs and are comparable in magnitude to that of the proton. The form factors of the $\Lambda$ and $\Sigma^0$, on the other hand, turn out to be very small. Predictions for the $\Sigma^-$ form factor are not reported in [36] but that form factor is presumably smaller in absolute value than the proton form factor. Assuming similar relative magnitudes of the form factors in the time-like region, we expect that, in the case of the hyperon channels listed in Tab. 3, the three-gluon contributions should match with the experimental data. This may not be the case for the $\Sigma^+\Sigma^+$ channel; while the three-gluon contribution is the same as for the $\Sigma^0\Sigma^0$ channel, the electromagnetic contribution may be large.

There is also a small contribution to the invariant $B^B_8$ from $c\bar{c}$ annihilations mediated by two gluons and a photon. The $gg\gamma$ contribution to $B^B_8$ being proportional to the $ggg$ contribution [33], amounts to less than about 1% of the latter and is therefore neglected.

Finally, from the measurement of the angular distribution of $B_8\bar{B}_8$ pairs produced in $e^+e^- \rightarrow J/\psi \rightarrow B_8\bar{B}_8$ [26] one observes small violations of the helicity sum rule: The fraction of $p\bar{p}$ and $\Lambda\bar{\Lambda}$ pairs with equal helicities amounts to about 10% of the total number of pairs with, in particular in the $\Lambda$ case, large errors. For the other hyperon channels the errors are so large that no conclusion can be drawn. The small amount of equal helicity pairs is what is to be expected if the process is dominated by perturbative QCD: Each of the virtual gluons creates a quark and an antiquark with opposite helicities. Since our baryon wave functions do not embody any non-zero orbital angular momentum component the quark helicities sum up to the baryon helicity. Hence, baryon and antibaryon are produced with opposite helicities. The small amount of $B_8\bar{B}_8$ pairs with the wrong helicity combination observed experimentally, while indicating the presence of some soft contributions, can be considered as a hint that perturbative QCD is the dominant dy-

\footnote{The time-like nucleon form factor is likely not under the regime of perturbative QCD in the energy region of interest; it is about a factor of 3 larger in absolute value than the form factor in the space-like region at $Q^2 = M_{\psi}^2$.}
ynamical mechanism in the $J/\psi \rightarrow B_8\overline{B}_8$ decays. One should, however, be aware of these contributions when theoretical results for the $J/\psi \rightarrow B_8\overline{B}_8$ decays are compared with experiment. The production of $B_8\overline{B}_8$ pairs with equal helicities can perhaps be explained as an constituent quark mass and/or baryon mass effect [33, 37].

5 $J/\psi$ decays into decuplet baryons

We are now going to apply the modified perturbative approach to the $J/\psi$ decays into the decuplet baryon-antibaryon channels $\Delta^+\Delta^-$ and $\Sigma^+\Sigma^-$ in the same manner as for the octet baryon case. As we said above, within a perturbative approach and with wave functions of zero orbital angular momentum components in the direction of the baryon momentum, $B_{10}$ and $B_{10}'$ are produced with opposite helicities. For the same reason only helicities $\pm 1/2$ are possible. Thus, as for the octet baryon case, the only non-zero helicity amplitudes are the $M_{\pm 1, \lambda}^{B_{10}}$. They are fed by only one invariant $B_{10}$ (or, depending on the definitions of the covariants, only one linear combination of the invariants) out of the five general covariant decomposition of the $J/\psi \rightarrow B_{10}\overline{B}_{10}$ helicity amplitudes comprises. Within the modified perturbative approach, the three-gluon contribution to the invariant $B_{10}$ reads

$$B_{3g}^{B_{10}} = \frac{3}{2} \bar{f} \int [dx][dx'] \int \frac{d^2b_1}{(4\pi)^2} \frac{d^2b_3}{(4\pi)^2} \hat{T}_H(x,x',b) \exp[-S(x,x',2m_c)]$$

$$\times \frac{1}{\lambda} \Psi_{123}(x,b)\Psi_{123}'(x',b), \quad (5.1)$$

for $B_{10} = \Delta, \Sigma^*$. The three-gluon contribution is evaluated with the decuplet wave functions introduced in Sect. 3 and with the hard scattering amplitude (4.7). Insertion of $B_{3g}^{B_{10}}$ into (4.2) provides the wanted decay widths. It still remains to choose plausible values of the parameters $a_{(10)}$ and $f_{(10)}$. The fact that we assume the same form for the $k_\perp$ dependence of the $\Delta$ and the nucleon wave functions, and that we use $\phi_{AS}$ for the $\Delta$ distribution amplitude which does not differ from the actual nucleon distribution amplitude (2.10) greatly, suggests, as a first attempt, the ansatz $a_{(10)} = a_{(8)} = 0.75 \text{ GeV}^{-1}$ and $P_{3q}^B = P_{3q}^N$. Thereby, it is perhaps plausible to evaluate the valence quark probability of the nucleon from $\phi_{AS}$ instead from (2.10). Doing so we find $P_{3q}^N(AS) = 0.163$ and the requirement $P_{3q}^B = P_{3q}^N(AS)$ leads to $f_{(10)} = 0.0163 \text{ GeV}^2$ which is larger than $f_{(8)}$ by a factor of $\sqrt{6}$, i.e., the SU(6) result. The results for the $J/\psi$ decay widths into $B_{10}\overline{B}_{10}$ pairs are presented in Tab. 4. As can be seen our result for the $\Sigma^*\Sigma^*$ channel is too large as compared to the data while agreement is achieved for the $\Delta\Delta$ case.

Bearing in mind that the $\Delta$ is in completely symmetric flavour and spin states and that the Pauli principle effectively induces an additional repulsive interquark force, one may expect a larger radius, and hence a larger value of $a_{(10)}$, for the $\Delta$ than for the nucleon. Therefore, we also try values for $a_{(10)}$ slightly larger than that for $a_{B_8}$ and fix in each case $f_{(10)}$ from the requirement $P_{3q}^B = 0.163$ as before. In Tab. 4 we list results obtained with the values $0.80$ and $0.85 \text{ GeV}^{-1}$ for $a_{(10)}$. For both these values of $a_{(10)}$ satisfactory agreement with experiment is found. The uncertainties of the theoretical results are the same as for the octet baryons. In particular, one has to consider the possibility of large electromagnetic contributions (see Fig. 2b).
<table>
<thead>
<tr>
<th>(a_{(10)}) [GeV(^{-1})]</th>
<th>(f_{(10)}) [GeV(^2)]</th>
<th>(B_1)</th>
<th>(B_2)</th>
<th>(B_3)</th>
<th>(B_4)</th>
<th>(B_5)</th>
<th>(\Gamma_{\Delta_{3g}}) [eV]</th>
<th>(\Gamma_{\Sigma^<em>\Sigma^</em>_{3g}}) [eV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.75</td>
<td>0.0163</td>
<td>-0.494</td>
<td>0.165</td>
<td>-0.203</td>
<td>-1.013</td>
<td>0.058</td>
<td>105</td>
<td>66.1</td>
</tr>
<tr>
<td>0.80</td>
<td>0.0143</td>
<td>-0.547</td>
<td>0.182</td>
<td>-0.216</td>
<td>-1.081</td>
<td>0.062</td>
<td>82.6</td>
<td>51.8</td>
</tr>
<tr>
<td>0.85</td>
<td>0.0127</td>
<td>-0.601</td>
<td>0.200</td>
<td>-0.229</td>
<td>-1.142</td>
<td>0.065</td>
<td>65.1</td>
<td>40.8</td>
</tr>
</tbody>
</table>

\(f_{(10)}\) is fixed by the requirement \(P_{\Delta_3} = 0.163\).

Table 4: \(f_{(10)}\), the expansion coefficients \(B_n\) for the \(\Sigma^*\) distribution amplitude and results for the three-gluon contribution to the \(J/\psi \to \Delta^{++}\Delta^{--}, \Sigma^*\Sigma^{**}\) decay widths for various values of the transverse size parameter \(a_{(10)}\) (\(m_s = 350\) MeV). \(f_{(10)}\) is fixed by the requirement \(P_{\Delta_3} = 0.163\).

### 6 Decays of the \(\psi'\) and other quarkonia

The extension of our approach to baryonic decays of the \(\psi'\) \((= \psi(2\,^3S_1))\) is now a simple matter. It is however important to realize that, in contrast to other authors [1], we evaluate the three-gluon contribution with the \(c\)-quark mass and not with the charmonium mass. This is, as we said, legitimate in a non-relativistic treatment of the charmonia. Hence, in order to get the \(\psi'\) widths in our approach we have not to rescale the corresponding \(J/\psi\) widths by \((M_{\psi}/M_{\psi'})^8\) but rather by

\[
\Gamma(\psi' \to BB) = \frac{\rho_{p.s.}(m_B/M_{\psi'})}{\rho_{p.s.}(m_B/M_{\psi})} \frac{\Gamma(\psi' \to e^+e^-)}{\Gamma(J/\psi \to e^+e^-)} \Gamma(J/\psi \to BB)
\]

which holds for both, octet and decuplet baryons. The charmonium decay constants are replaced by the leptonic decay widths by means of formula (4.11).

As an immediate examination of our approach one may apply the scaling relation (6.1) directly to the experimental data (see Tabs. 3, 4, 5). Considering the uncertainties due to the electromagnetic contributions, which in one or the other case may be large, (6.1) works quite well in particular for the octet baryons. For the decuplet baryons, on the other hand, it seems that the suppression of the \(\psi'\) decay widths is a slightly underestimated, although the large experimental errors prevent any definite conclusion at present. One may suspect the neglect of the decuplet baryon masses to be responsible for that possible imperfection. In any case, an additional strong suppression, as provided by \((M_{\psi}/M_{\psi'})^8\) [1], is in clear conflict with the data [9]. This observation supports our attempt of using the \(c\)-quark mass in the calculation of the decay amplitudes rather than the mass of the charmonium state in question.

Results for baryonic decay widths of the \(\psi'\), evaluated through (6.1) from the set 3 \(J/\psi\) widths, are listed in Tab. 5 where also recent experimental results of the BES collaboration [9] are quoted. The data are still preliminary. The agreement between theoretical results (with \(a_{(10)} = 0.85\) GeV\(^{-1}\) in the decuplet baryon case) and experiment is generally good although our results seem to be a bit too large for the \(\Sigma\Sigma\) and \(\Sigma^*\Sigma^*\) channels. For a discussion of uncertainties we refer to Sect. 3.

Computation of the \(\psi(3\,^3S_1)\) decay widths are difficult within our approach. The relativistic corrections are presumably larger since the \(\psi(3\,^3S_1)\) mass is above the threshold for open charm production. The \(BB\) decay widths are likely to be very small and it is
Table 5: The three-gluon contribution to the $\psi' \to B\bar{B}$ decay widths (in [eV]) computed through (6.1) from the set 3 $J/\psi$ widths (see Tab. 3). $a_{(10)} = 0.85$ GeV$^{-1}$.

Results for bottomonium decays, on the other hand, can safely be calculated within our approach. The hard scale, provided by the $b$-quark mass, is larger than in the charmonium case and relativistic corrections are smaller. But, as it turns out, the predicted decay widths for the baryonic channels are also very small. Approximately, i.e. ignoring the fact that the $k_\perp$-dependent suppression of the three-gluon contribution is a little bit different in the two cases, we find the following rescaling formula

$$\Gamma(\Upsilon \to B\bar{B}) = 4 \rho_{p.s.}(m_B/M_\Upsilon) \frac{\Gamma(\Upsilon \to e^+e^-)}{\rho_{p.s.}(m_B/M_\psi) \Gamma(J/\psi \to e^+e^-)} \left(\frac{m_c}{m_b}\right)^8 \Gamma(J/\psi \to B\bar{B})$$

Using $m_b = 4.5$ GeV we obtain, for instance, a value of 0.03 eV for the $\Upsilon \to p\bar{p}$ decay width which value corresponds to a branching ratio of $0.5 \times 10^{-6}$ well below the experimental upper bound [24]. The decay widths for the other $B\bar{B}$ channels are even smaller.

## 7 Summary

In this investigation, we applied the modified perturbative approach to the decays of $J/\psi$ and $\psi'$ into baryon-antibaryon pairs. We demonstrated that, on the basis of plausible baryon wave functions for which SU(3)$_F$ symmetry is only mildly broken by quark mass effects and for which even SU(6) symmetry, in sharp contrast to the QCD sum rule based wave functions, approximately holds, the experimental data for octet and decuplet baryon channels are quite well reproduced by the phase space corrected three-gluon contributions. The perturbative contributions to the decay widths are calculated self-consistently in the sense that the bulk of a perturbative contribution is accumulated in regions of reasonably small values of $\alpha_s$. Besides the form of the wave functions used by us our analysis differs from previous ones in the following points:

i) The use of the modified perturbative approach allows to take into account the running $\alpha_s$ and the evolution of the wave functions properly, in contrast to the usual leading twist analysis. The virtualities of the internal $c$ quarks and gluons can be chosen as the arguments of $\alpha_s$. Hence, $\alpha_s$ reflects the characteristic scale of the process under question. The running coupling constant is therefore not a quasi-free parameter that can, within a certain range, be chosen arbitrarily. Since the decay widths are proportional to $\alpha_s^6$, a large factor of uncertainty is therefore hidden in the standard perturbative analysis.

ii) The hard scattering amplitude is computed with the $c$-quark mass instead of the
charmonium mass. The use of the $c$-quark mass is consistent with the non-relativistic treatment of the charmonium state and with the perturbative approach. Hence the $\psi'$ - $J/\psi$ scaling relation (6.1) holds in our approach. That scaling relation is nicely confirmed by the data in contrast to the usual $(M_\psi/M_{\psi'})^8$ scaling.

Although our results for the decay widths agree with the data, we are aware of a number of uncertainties in our calculation. Above all, we mention as a source of uncertainty the value of the $c$-quark mass. The value we have chosen (1.5 GeV) is consistent with other constraints on that mass. Another uncertainty arises from the electromagnetic contribution which, at least in the $p\bar{p}$ case, can be large. In virtue of the unknown relative phase between the electromagnetic and the three-gluon contribution, the first cannot be taken into account properly. We emphasize that most of the uncertainties cancel in ratios of widths to a large extent.

It should be noted that we have not found hints at substantial contributions from the $c\bar{c}g$ Fock state, i.e., colour octet contributions [18], neither in the baryonic $J/\psi$ nor in the $\psi'$ decays.

An interesting class of $J/\psi$ decays are the $B_8\overline{B}_{10}$ channels. While the three-gluon contributions to the $p\Delta^-$ and $\Sigma^{*0}\Lambda$ channels are strictly zero [38] they are non-zero - although small - for other $B_8\overline{B}_{10}$ channels, since our wave functions exhibit only a mild breaking of SU(3)$_F$ symmetry. Experimentally, only upper bounds are known for the first two channels [24] saying that these decays are indeed suppressed by at least an order of magnitude as compared to, say, the $p\bar{p}$ or the $\Lambda\bar{\Lambda}$ channels. The experimental widths for the other $B_8\overline{B}_{10}$ channels are surprisingly large [24]. Thus, in accordance with [38], we expect as the dominant SU(3)$_F$ breaking mechanism for these reactions a sizeable electromagnetic contribution.

Finally, we have to mention that there are a few exclusive charmonium decays which cannot be described within the standard or the modified perturbative approach. Thus, the relatively large branching ratio of the process $J/\psi \rightarrow \rho\pi$ observed experimentally, for instance, indicates a substantial violation of hadronic helicity conservation while only mild violations are observed in the baryonic $J/\psi$ decays (see the discussion in Sect. 3). For a discussion of this puzzle and a possible solution of it by means of a hypothetical glueball see [39].

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