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Exclusive Decuplet-Baryon Pair Production in Two-Photon Collisions

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ABSTRACT

This work extends our previous studies of two-photon annihilation into baryon-antibaryon pairs from spin-1/2 octet to spin-3/2 decuplet baryons. Our approach is based on perturbative QCD and treats baryons as quark-diquark systems. Using the same model parameters as in our previous work, supplemented by QCD sum-rule results for decuplet baryon wave functions, we are able to give absolute predictions for decuplet baryon cross sections without introducing new parameters. We find that the Δ^{++} cross section is of the same order of magnitude as the proton cross section, well within experimental bounds.

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1 Introduction

The study of exclusive processes in quantum chromodynamics (QCD), where intact hadrons are explicitly measured in the final state, provides important insights into the mechanisms of confinement and into the dynamics of hadronic bound states [1, 2]. Among the multitude of exclusive processes, two-photon annihilation into baryon-antibaryon pairs is particularly interesting, because it is one of the simplest calculable large-angle hadronic scattering reactions involving two hadrons. Therefore, $\gamma\gamma \to \mathrm{B}\bar{\mathrm{B}}$ has recently received considerable experimental [3] and theoretical [4, 5, 6] attention.

In a recent paper [5] we have studied baryon pair production in two-photon collisions for baryons belonging to the lowest-lying flavor octet. In the present note we extend our work to reactions involving spin-3/2 decuplet baryons. Previously, two-photon annihilation into decuplet baryons has been studied in Refs. [7, 8, 9, 6, 10] within different frameworks with differing conclusions. Thus an experimental analysis could shed light on the relative importance of the underlying mechanisms considered here and in the aforementioned references.

Our model is a modification of the perturbative hard-scattering picture (HSP) for exclusive processes [11, 12]. While the HSP is exactly valid only at asymptotically large momentum transfer, the interplay of perturbatively calculable with nonperturbative effects renders theoretical analyses quite intricate at energies where data are currently available. In order to parameterize such possible non-perturbative effects within a perturbative framework, an effective formalism was developed in Ref. [13], where baryons are treated as quark-diquark systems. In the sequel this model has been successfully applied to a variety of exclusive reactions [5, 8, 14, 15, 16, 17].

In the following, we start with a brief review of the quark-diquark model. Then we go on to describe the new ingredients necessary for the study of processes involving decuplet baryons. In Sec. 3 we present and discuss model predictions with emphasis on the Δ cross sections, for which experimental upper bounds are available [18]. Following concluding remarks, supplementary analytical expressions for the scattering amplitudes are tabulated in the Appendix.

2 Exclusive Reactions in the Quark-Diquark Picture

Here we briefly summarize the modified hard-scattering formalism with diquarks, and elaborate on the aspects specific to the treatment of decuplet baryons. For a full account of all details we refer to our recent work [5, 17].

2.1 Review of the Model

As in the conventional hard-scattering picture [11, 12], an exclusive reaction amplitude $\mathcal{M}$ is convolutively factorized into a process-dependent, perturbative hard-scattering amplitude $\hat{T}$ and process-independent, non-perturbative distribution amplitudes Ψ. The latter are probability amplitudes for finding the pertinent valence Fock states, here quarks and diquarks, in the scattering hadrons. The amplitude for two-photon annihilation into a baryon-antibaryon pair is given by

$$\overline{\mathcal{M}}_{\{\lambda\}}\left(\hat{s},\hat{t}\right) \; = \int_0^1 dx_1 \int_0^1 dy_1 \Psi^\dagger_{\mathrm{B}}\left(x_1\right) \Psi^\dagger_{\overline{\mathrm{B}}}\left(y_1\right) \hat{T}_{\{\lambda\}}\left(x_1, y_1; \hat{s}, \hat{t}\right), \tag{1}$$

where Lorentz and color indices are suppressed for convenience. Furthermore, the dependence on renormalization and factorization scales is neglected since we are only interested in a rather restricted range of momentum transfer. The subscript $\{\lambda\}$ denotes all possible configurations of photon and baryon helicities. In the following we use the label B to denote spin-1/2 octet baryons and B_{10} to label spin-3/2 decuplet baryons.

For the process $\gamma\gamma \to \mathrm{B}_{10}\overline{\mathrm{B}}_{10}$, there are 19 independent helicity amplitudes, $\overline{\mathcal{M}}_{\lambda_{\mathrm{B}_{10}},\lambda_{\overline{\mathrm{B}}_{10}};\lambda_1,\lambda_2}$, where the $\lambda_{\mathrm{B}_{10}}$, $\lambda_{\overline{\mathrm{B}}_{10}}$ are the helicities of the outgoing baryon and antibaryon, respectively, and λ_1, λ_2 label the helicities of the two photons. Only 13 out of these 19 helicity amplitudes involve a zero or single flip of the hadronic helicity. Double flip amplitudes vanish in our approach. We use the following convention for the nonvanishing amplitudes:

$$\begin{array}{ll}
\overline{\phi}_1 = \overline{\mathcal{M}}_{-\frac{1}{2},\frac{1}{2};1,-1}, & \overline{\phi}_7 = \overline{\mathcal{M}}_{-\frac{1}{2},\frac{3}{2};1,-1}, \\
\overline{\phi}_2 = \overline{\mathcal{M}}_{-\frac{1}{2},-\frac{1}{2};1,1}, & \overline{\phi}_8 = \overline{\mathcal{M}}_{\frac{1}{2},-\frac{3}{2};1,1}, \\
\overline{\phi}_3 = \overline{\mathcal{M}}_{\frac{1}{2},-\frac{1}{2};1,1}, & \overline{\phi}_9 = \overline{\mathcal{M}}_{\frac{1}{2},-\frac{3}{2};1,-1}, \\
\overline{\phi}_4 = \overline{\mathcal{M}}_{\frac{1}{2},\frac{1}{2};1,-1}, & \overline{\phi}_{10} = \overline{\mathcal{M}}_{-\frac{1}{2},\frac{3}{2};1,1}, \\
\overline{\phi}_5 = \overline{\mathcal{M}}_{\frac{1}{2},-\frac{1}{2};1,-1}, & \overline{\phi}_{11} = \overline{\mathcal{M}}_{-\frac{3}{2},\frac{3}{2};1,-1}, \\
\overline{\phi}_6 = \overline{\mathcal{M}}_{\frac{1}{2},\frac{1}{2};1,1}, & \overline{\phi}_{12} = \overline{\mathcal{M}}_{\frac{3}{2},-\frac{3}{2};1,1}, \\
 & \overline{\phi}_{13} = \overline{\mathcal{M}}_{\frac{3}{2},-\frac{3}{2};1,-1}.
\end{array} \tag{2}$$

Other helicity configurations are related to these via parity and/or time reversal invariance. Our normalization of the amplitudes is such that the differential cross section for two-photon annihilation into decuplet baryons is given by

$$\frac{d\sigma}{dt} = \frac{1}{64\pi s^2} \sum_{\{\lambda\}} \left|\overline{\mathcal{M}}_{\{\lambda\}}\right|^2, \tag{3}$$

where the sum is over all possible helicity configurations $\{\lambda\}$.

In (1), $\hat{T}$ consists of all possible tree diagrams that contribute to the elementary scattering process $\gamma\gamma \to qD\bar{q}\bar{D}$. The momenta carried by quarks q and diquarks D are assumed to be collinear to those of their parent hadrons, B. The quark and antiquark carry momentum fractions x_1 and y_1 in the baryon and antibaryon, respectively, while the diquark and antidiquark carry momentum fractions $x_2 = 1 - x_1$ and $y_2 = 1 - y_1$, respectively. Since we assume that every baryonic constituent has a four-momentum $x\, p_{\rm B}$ proportional to the four-momentum of its parent hadron $p_{\rm B}$ [19], it acquires an effective mass $xm_{\rm B}$, where $m_{\rm B}$ denotes the baryon mass. These effective masses are taken into account for all internal and external legs of the Feynman diagrams contributing to the hard-scattering amplitude $\hat{T}$. The hard-scattering amplitude is then expanded in powers of the small parameter $(m_{\rm B}/\sqrt{s})$ up to next-to-leading order, at fixed center-of-mass scattering angle $\hat{\theta}$. The result is reexpressed in terms of massless Mandelstam variables, $\hat{s}$, $\hat{t}$, and $\hat{u}$ which are obtained from the usual massive Mandelstam variables, s, t, u, again by expansion in $(m_{\rm B}/\sqrt{\hat{s}})$. In the hard scattering diagrams, the composite nature of the diquarks is taken into account by diquark form factors. These are parameterized such that asymptotically the scaling behavior of the pure quark HSP emerges.

The complete parameterization of the model, including form factors and octet-baryon wave functions can be found in [5]. These parameters were fixed in [14] by fitting elastic electron-nucleon scattering data. With the same set of parameters a variety of other processes has been computed, and the results have successfully met experimental comparison [5, 14, 16, 17].

2.2 Decuplet Baryons

The diquark model comprises spin-0 (scalar) and spin-1 (vector) diquarks. While both scalar (S) and vector (V) diquarks contribute to processes involving spin-1/2 octet baryons, the valence Fock states of spin-3/2 decuplet baryons consist only of quarks and vector diquarks.

We recall that the valence Fock state of an octet baryon B with mass $m_{\rm B}$, momentum $p_{\rm B}$, and helicity λ can be described by the following quark-diquark wave function

$$\Psi_{\rm B}(p_{\rm B}, x, \lambda) = f_S^{\rm B} \Phi_S^{\rm B}(x) \chi_S^{\rm B}\, u(p_{\rm B}, \lambda) + f_V^{\rm B} \Phi_V^{\rm B}(x) \chi_V^{\rm B} \frac{1}{\sqrt{3}} \left(\gamma^\mu + \frac{p_{\rm B}^\mu}{m_{\rm B}} \right) \gamma_5\, u(p_{\rm B}, \lambda) \qquad (4)$$

when transverse momenta of the constituents are neglected. x is the longitudinal momentum fraction of the quark, whereas the diquark carries the longitudinal momentum fraction $1 - x$. Analogously, the wave function of a decuplet baryon may be written as

$$\Psi^{\mu}_{\mathrm{B}_{10}}(p_{\mathrm{B}_{10}}, x, \lambda) = f_V^{\mathrm{B}_{10}} \Phi_V^{\mathrm{B}_{10}}(x) \chi_V^{\mathrm{B}_{10}} u^{\mu}\left(p_{\mathrm{B}_{10}}, \lambda\right), \tag{5}$$

with the Rarita-Schwinger spinors [20]

$$\begin{aligned} u^{\mu}\left(p, \lambda = \pm 3/2\right) &= \varepsilon^{\mu}\left(\pm 1\right)\, u\left(p, \lambda = \pm 1/2\right)\,, \\ u^{\mu}\left(p, \lambda = \pm 1/2\right) &= \left[\sqrt{\frac{3}{2}}\varepsilon^{\mu}(0) - \frac{2\lambda}{\sqrt{6}}\left(\gamma^{\mu} + \frac{p^{\mu}}{m_{\mathrm{B}_{10}}}\right)\gamma_5\right] u(p, \lambda). \end{aligned} \tag{6}$$

Recall that all Lorentz indices have been suppressed in Eq. (1), the open index μ of the vector diquark polarization vector in (4) and (5) is contracted appropriately in the convolution integral (1). χ_D^{B}, $\chi_D^{\mathrm{B}_{10}}$ $(D = S, V)$ denote pertinent SU(3) quark-diquark flavor wave functions and Φ_D^{B}, $\Phi_D^{\mathrm{B}_{10}}$ represent the nonperturbative probability amplitudes for finding these constituents with momentum fractions x and $1 - x$, respectively, in the (decuplet) baryon. These probability amplitudes are normalized such that

$$\int_0^1 dx \Phi_D^{\mathrm{B}}(x) = 1, \tag{7}$$

and analogously for $\Phi_D^{\mathrm{B}_{10}}$. The constants f_D^{B}, $f_D^{\mathrm{B}_{10}}$ result from integrating out intrinsic transverse momenta in the full wave function to produce Eqs. (4) and (5), respectively. The numerical values of f_D^{B} and $f_D^{\mathrm{B}_{10}}$ are furthermore determined by the overall probability of finding the $|qD\rangle$-state in the baryon B or decuplet baryon B_{10}, respectively.

For unbroken SU(6) spin-flavor symmetry octet- and decuplet baryon wave functions are related, specifically, $\Phi_S^{\mathrm{B}} = \Phi_V^{\mathrm{B}} = \Phi_V^{\mathrm{B}_{10}}$ and $f_S^{\mathrm{B}} = f_V^{\mathrm{B}} = f_V^{\mathrm{B}_{10}}/\sqrt{2}$. In the actual parameterization of the diquark model [14] the asymptotic SU(6) symmetry is systematically broken down to SU(3) flavor symmetry. Thus the above SU(6) relations are by no means satisfied, and Φ_S^{B} and Φ_V^{B} as well as f_S^{B} and f_V^{B} have quite different values. Since SU(6) symmetry is thus already broken within the baryon octet we cannot use SU(6) symmetry for deriving quark-diquark wave functions of decuplet baryons. Instead, we will apply another strategy to fix $\Phi_V^{\mathrm{B}_{10}}$ and $f_V^{\mathrm{B}_{10}}$.

The lowest moments of three-quark wave functions of octet and decuplet baryons are restricted by QCD sum rules [9, 21]. Model wave functions that satisfy the QCD sum-rule constraints (for a typical factorization scale of about 1 GeV) are very asymmetric in the longitudinal momentum fractions x_i, $i = 1, 2, 3$ of the quarks for octet baryons and nearly symmetric ($\sim x_1 x_2 x_3$) for the Δs and the Ω [9]. By regrouping terms in the three-quark wave function such that, for example, quarks 2 and 3 are in a specific spin-flavor state and by integrating over one of the momentum fractions of the two quarks that build up this "diquark" we can convert the three-quark wave function into a quark-diquark wave function that nearly has the form

(4) or (5) for octet or decuplet baryons, respectively. For more information on this conversion we refer to [8]. The probability amplitudes Φ_V^{B} and $\Phi_V^{\mathrm{B}_{10}}$ for general three-quark wave functions are different in the cases of helicity-0 and helicity-1 V diquarks. For the octet and decuplet model wave functions that we employ this difference turns out to negligible. We then arrive at Eq. (4) or (5), respectively.

We apply the above procedure to the three-quark wave function of the Δ that has been proposed in Ref. [9] based on QCD sum-rule constraints. We obtain the following quark-diquark wave function for a Δ with helicity $\pm 1/2$

$$\Phi_V^{\Delta,|\lambda|=1/2}(x) = Nx(1-x)^3(1-2.95x+3.86x^2)\exp\left\{-b^2\left[\frac{m_q^2}{x}+\frac{m_V^2}{1-x}\right]\right\}. \tag{8}$$

Analogous to the standard parameterization of the diquark model for octet baryons [14] we have introduced an additional exponential factor that damps the end-point regions $x \to 0, 1$. Such an exponential factor results if the transverse momentum dependence of the full wave function, which is integrated over, is assumed to be of Gaussian form. The parameters $b^2 = 0.248$ GeV^2, $m_q = 0.33$ GeV, and $m_V = 0.58$ GeV are taken to be the same as for octet baryons. The normalization factor N is determined by Eq. (7). The expression for $\Phi_V^{\Delta,|\lambda|=3/2}(x)$ differs, in general, from $\Phi_V^{\Delta,|\lambda|=1/2}(x)$. However, we refrain from quoting it here, because our explicit calculations show that the production of helicity-3/2 Δs is suppressed within the diquark model.

The only remaining open parameter is now the normalization $f_V^{\Delta,|\lambda|=1/2}$ of the helicity-1/2 Δ wave function. Since the normalization f_V^{B} of the octet-baryon wave function was taken as a free parameter in the diquark model we normalize the Δ wave function relative to the proton wave function. This means that we convert the three-quark wave functions for proton and Δ into quark-diquark wave functions of the form (4) and (5), respectively, and consider the resulting ratio $f_V^{\Delta,|\lambda|=1/2}/f_V^{\mathrm{p}}$. For the QCD sum-rule based wave functions of Refs. [9] and [21] this ratio becomes $f_V^{\Delta,|\lambda|=1/2}/f_V^{\mathrm{p}} = 0.898$. With $f_V^{\mathrm{p}} = 127.7$ MeV, the value obtained in a fit of elastic electron-nucleon scattering data [14], we thus find

$$f_V^{\Delta,\,|\lambda|=1/2} = 125.1 \text{ MeV}\,. \tag{9}$$

This completes the parameterization of our model for decuplet baryons. For sake of completeness we quote the flavor wave functions entering (5) for the differently charged Δs:

$$\begin{aligned}
\chi_V^{\Delta^{++}} &= uV_{\{uu\}}\,, \\
\chi_V^{\Delta^{+}} &= \left[\sqrt{2}uV_{\{ud\}} + dV_{\{uu\}}\right]/\sqrt{3}\,, \\
\chi_V^{\Delta^{0}} &= \left[\sqrt{2}dV_{\{ud\}} + uV_{\{dd\}}\right]/\sqrt{3}\,, \\
\chi_V^{\Delta^{-}} &= dV_{\{dd\}}\,.
\end{aligned} \tag{10}$$

3 Results

We list analytical results for the hard-scattering amplitudes $\hat{T}_{\{\lambda\}}$ contributing to $\gamma\gamma \to \mathrm{B}_{10}\bar{\mathrm{B}}_{10}$ in the Appendix. These results have been checked via crossing relations [22] against the separately computed amplitudes for the crossed process, Compton scattering $\gamma\mathrm{B}_{10} \to \gamma\mathrm{B}_{10}$. Comparing the spinor structure of the decuplet baryon wave function (5) with the one for octet baryons (4), we find that the leading, non-flip, hard amplitudes for decuplet baryons with helicity 1/2 are related by a factor of 2 to those for octet baryons. From the analytical expressions we also observe, that the hard-scattering amplitudes for octet baryons with helicity $\pm 3/2$ are suppressed by $\mathcal{O}(m^2_{\mathrm{B}_{10}}/\hat{s})$ or higher, even if these amplitudes conserve the hadronic helicity or flip it by one unit. The only 4-point contribution that is not suppressed enters the helicity amplitude $\bar{\phi}_2$. In the numerical calculations this contribution, however, turns out to be nearly negligible. 5-point functions with both photons attaching to the diquark do not contribute at all, since these are also suppressed by $\mathcal{O}(m^2_{\mathrm{B}_{10}}/\hat{s})$ or even higher.

As a consequence of this observation, cross section ratios of different decuplet baryon channels can easily be estimated, provided that the corresponding probability amplitudes $\Phi_V^{\mathrm{B}_{10}}$ are not too different. The cross-section ratios are then essentially determined by the corresponding charge-flavor factors $C_{\mathrm{cf}}^{(3)}$ (see Eq.(13)) and the wave function normalizations $f_V^{\mathrm{B}_{10}}$. For the Δ-quartet Φ_V^{Δ} and f_V^{Δ} are the same for all members due to isospin symmetry. From the flavor wave functions (10) the charge-flavor factors $C_{\mathrm{cf}}^{(3)}$ are seen to be 4/9, 3/9, 2/9, and 1/9 for the Δ^{++}, Δ^+, Δ^0, and Δ^-, respectively. The cross section ratios become (approximately)

$$\sigma(\Delta^{++}) : \sigma(\Delta^{+}) : \sigma(\Delta^{0}) : \sigma(\Delta^{-}) = 16 : 9 : 4 : 1\,. \tag{11}$$

This is the first interesting prediction of the diquark model. In Fig. 1 we show the integrated cross sections ($|\cos(\theta_{CM})| < 0.6$, where θ_{CM} is the center-of-mass scattering angle) for the Δ channels. The plot exhibits numerical predictions obtained with the standard parameterization of the diquark model [5] and the Δ wave function derived in Sec. 2.2. It confirms Eq. (11) within 1 percent.

This prediction is to be contrasted with the ratios 16 : 1 : 0 : 1 that result if the photons couple to the total charge of the Δs. Also within the pure quark HSP the ratios for the Δ^+ and the Δ^0 channels differ from ours. Within the pure quark HSP the cross section ratios for the different Δ channels are predicted to be $\sigma(\Delta^{++}) : \sigma(\Delta^{+}) : \sigma(\Delta^{0}) : \sigma(\Delta^{-}) \approx 16 : 2 : 1/3 : 1$ [9]. Note that all the above predictions agree with our result for the cross section ratio $\sigma(\Delta^{++}) : \sigma(\Delta^{-}) \approx 16 : 1$. This result is also found in a more general QCD analysis [10]. However, yet another possible production mechanism via multi-pion intermediate states predicts $\sigma(\Delta^{++}) = \sigma(\Delta^{-})$ and $\sigma(\Delta^{+}) = \sigma(\Delta^{0})$ [6]. An experimental determination of such cross section ratios

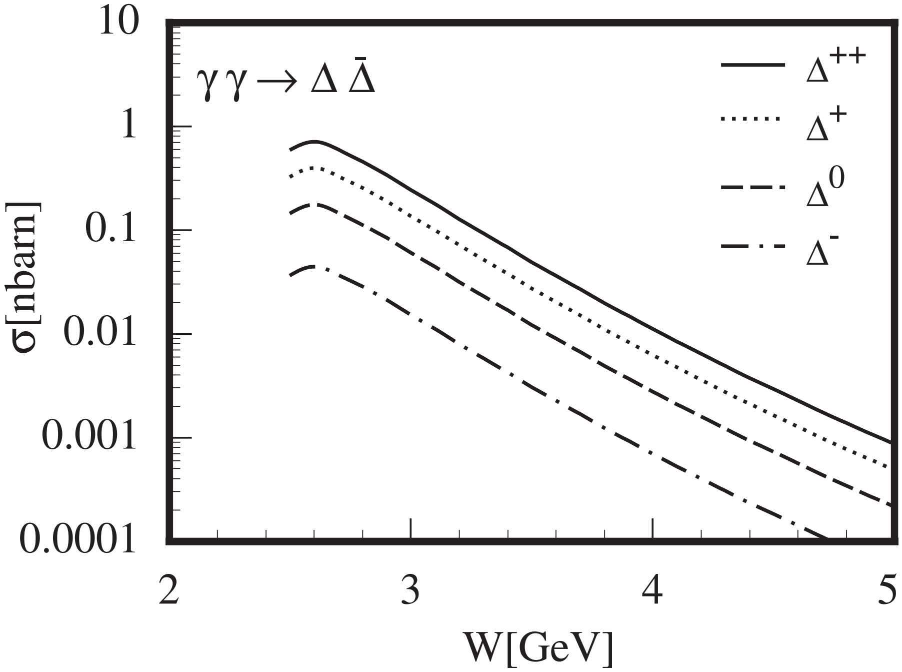

Figure 1: Integrated cross sections for $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ (solid line), $\Delta^{+}\bar{\Delta}^{-}$ (dotted), $\Delta^{0}\bar{\Delta}^{0}$ (dashed), $\Delta^{-}\bar{\Delta}^{+}$ (dash-dotted line) ($|\cos(\theta_{CM})| < 0.6$) versus center-of-mass energy $W = \sqrt{s}$ predicted with the standard parameterization of the diquark model [5] and the Δ DA defined in the text (see Eqs. (5), (8), and (9)).

could therefore provide important clues on the underlying production mechanisms, especially because in ratios of cross sections for different Δ channels the sensitivity to the specific form of the Δ wave function should be greatly reduced.

If we assume SU(3)-flavor symmetry, that is, if we take the same $\Phi_V^{\mathrm{B}_{10}}$ and $f_V^{\mathrm{B}_{10}}$ for all decuplet baryons, we are also able to give estimates for the pair production of strange decuplet baryons. Aside from appropriate phase space factors, SU(3) symmetry implies

$$\begin{aligned}
\sigma(\Delta^{+}) &= \sigma(\Sigma^{*+})\,, \\
\sigma(\Delta^{0}) &= \sigma(\Sigma^{*0}) = \sigma(\Xi^{*0})\,, \\
\sigma(\Delta^{-}) &= \sigma(\Sigma^{*-}) = \sigma(\Xi^{*-}) = \sigma(\Omega^{*-})\,.
\end{aligned} \tag{12}$$

However, since it is experimentally very difficult to measure pair-production cross sections for decuplet baryons, we refrain from giving quantitative results for the strange decuplet baryons. We rather concentrate in the following on the Δ^{++} channel which might have the best chance to be measured due to its comparably large cross section.

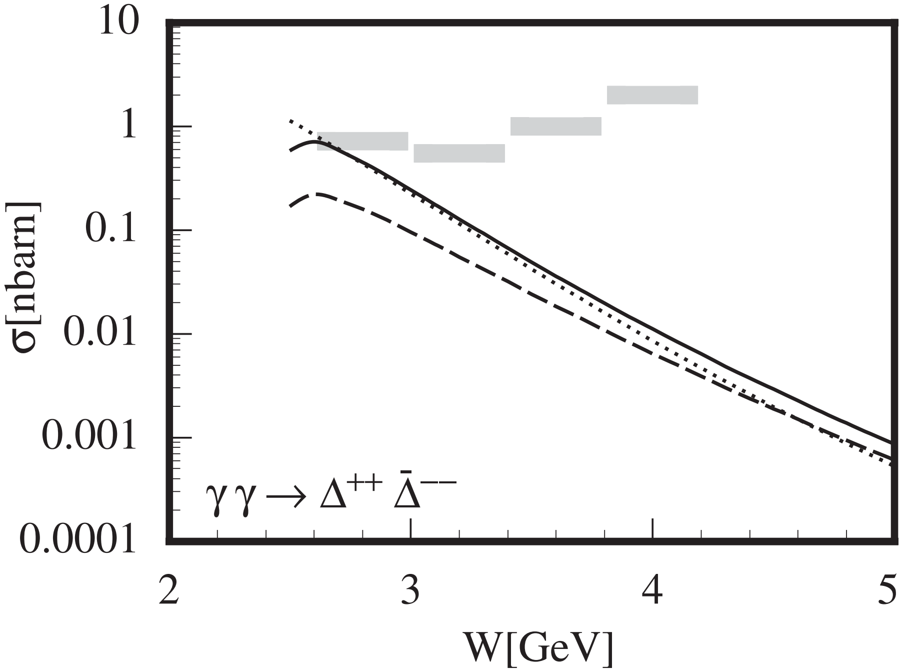

Figure 2: Integrated cross section for $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ ($|\cos(\theta_{CM})| < 0.6$) versus $W = \sqrt{s}$ for the same Δ DA as in Fig 2. The solid line corresponds to the full diquark-model calculation. The contribution to the cross section that comes from the hadronic-helicity conserving amplitudes $\bar{\phi}_1$ and $\bar{\phi}_5$ is represented by the dashed line. For comparison we also show the integrated cross section for $\gamma\gamma \to \mathrm{p}\bar{\mathrm{p}}$ (dotted line) calculated within the same model [5]. The shaded boxes indicate experimental upper bounds as obtained by the ARGUS collaboration [18].

In Fig. 2 we show for comparison with the $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ cross section the $\gamma\gamma \to \mathrm{p}\bar{\mathrm{p}}$ cross section that we have obtained with the same parameterization [5]. Surprisingly, we find that the Δ^{++} cross section is of the same order of magnitude as the proton cross section. This prediction seems to be very stable against (reasonable) changes of the Δ wave function. With a Δ wave function that satisfies the SU(6) relations $\Phi_V^\Delta = \Phi_V^\mathrm{p}$ and $f_V^\Delta = \sqrt{2} f_V^\mathrm{p}$ we obtain, for example, a result which is only about 20% to 30% smaller *. Under the naive assumption that the photons couple directly to the charges of the baryons one would expect the Δ^{++} cross section to be about 16 times larger than the proton cross section.

*In a previous attempt to estimate $\sigma(\Delta^{++})/\sigma(\mathrm{p})$ within a diquark model a ratio of ≈ 0.1 was found [8]. This, however, was obtained with an incomplete version of the diquark model, where V diquarks were not taken fully into account and mass effects have been neglected.

From the viewpoint of the pure quark hard-scattering picture, the ratio of $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ to the $\gamma\gamma \to \mathrm{p}\bar{\mathrm{p}}$ cross section depends strongly on the choice of the proton wave function [9]. Not surprisingly, a result for the ratio comparable to ours is obtained with the QCD sum-rule wave functions of Refs. [9] and [21] for Δ and proton, respectively, which we have used in Sec. 2.2 to derive and normalize our quark-diquark wave function of the Δ. However, if the asymptotic wave function $\sim x_1x_2x_3$ is taken for both the proton and the Δ, the cross section ratio $\sigma(\Delta^{++})/\sigma(\mathrm{p})$ can be as large as 50 within the pure quark HSP [7]. On the other hand, soliton models involving multi-pion channels predict a much smaller ratio [2, 6], comparable to our findings. An experimental determination of the ratio $\sigma(\Delta^{++})/\sigma(\mathrm{p})$ could therefore help to explore the importance of the various mechanisms that result in these quite different predictions.

Unfortunately, it is very difficult experimentally to isolate the signal of the broad Δ^{++} resonance from the background and to disentangle the Δ^{++} and the Δ^0 contributions in the $\gamma\gamma \to \mathrm{p}\bar{\mathrm{p}}\pi^+\pi^-$ cross sections which are actually measured. Therefore only upper limits for the $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ cross section have been extracted up to now by the ARGUS collaboration [18]. As can be seen in Fig. 2, our results lie well below these upper limits. More recent attempts to constrain the $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ cross section using the data taken by the L3 group are afflicted with the same problems, but a preliminary assessment indicates compatibility with the ARGUS results and our predictions [23]. A better chance to determine the cross section for $\Delta^{++}\bar{\Delta}^{--}$ pair production would perhaps exist for the BABAR or BELLE experiments which enjoy a much higher luminosity.

Finally, let us comment on the treatment of mass effects within our approach. Fig. 2 displays the effect of taking into account the finite Δ mass. As explained in Sec. 2.1, the Δ mass is taken into account in the hard-scattering amplitudes via an expansion in the small parameter $(m_{\mathrm{B}}/\sqrt{\hat{s}})$ where only the leading and next-to leading order terms are kept. As expected, mass correction terms do not contribute to the hadronic helicity-conserving amplitudes $\overline{\phi}_1$ and $\overline{\phi}_5$. Only the amplitudes that involve a single flip of the hadronic helicity, which vanish if masses are neglected, become nonzero due to the mass correction terms. The comparison of the solid and the dashed lines in Fig. 2 shows that these mass effects can be sizable in the few-GeV region. At $W = 2.5$ GeV the leading-order contributions provide only about 30% of the full cross section. This ratio increases, of course, with increasing energy and becomes roughly 70% at $W = 5$ GeV.

4 Concluding Remarks

In this work we have computed $\gamma\gamma \to \mathrm{B}_{10}\bar{\mathrm{B}}_{10}$ cross sections at intermediate momentum transfer for the case of spin-3/2 decuplet baryons B_{10}. We have employed a modification of the hard-scattering picture for exclusive reactions, where baryons are treated as quark-diquark systems, thereby effectively parameterizing nonperturbative contributions which are undoubtedly present at currently experimentally accessible energies. Using the same model parameters as in previous studies of other photon-induced reactions, and constraining the quark-diquark wave function of the Δ with the help of QCD sum-rule results, we are able to give absolute predictions for $\gamma\gamma \to \Delta\bar{\Delta}$ without introducing new parameters.

We find that the cross section for $\gamma\gamma \to \Delta^{++}\bar{\Delta}^{--}$ is of the same order of magnitude as the cross section for proton pair production, $\gamma\gamma \to p\bar{p}$. Furthermore, we observe that the pair production of decuplet baryons is almost completely determined within our model by those graphs where both photons couple to the quark line. This enables us to estimate production ratios for different decuplet-baryon channels independent of the choice of the wave function, provided that the wave functions are similar for all baryons within the decuplet. This is certainly the case for the Δ-quartet for which we predict the ratios $\sigma(\Delta^{++}) : \sigma(\Delta^{+}) : \sigma(\Delta^{0}) : \sigma(\Delta^{-}) = 16 : 9 : 4 : 1$.

There are various other estimates of these cross section ratios in the literature, based on different viewpoints and production mechanisms, which differ in their predictions from ours. It would therefore be necessary to compare to experimental analyses, in order to determine the relative importance of the considered production mechanisms, and to learn more about the degree of symmetry among constituents in decuplet-baryon distribution amplitudes. Such an experimental analysis should be quite feasible at a high-luminosity e^+e^- collider. We therefore hope that our experimental colleagues will study this interesting problem in the near future.

A Elementary Helicity Amplitudes for $\gamma\gamma \to qV\bar{q}\bar{V}$

There are 30 Feynman graphs that contribute to the hard-scattering amplitudes $\hat{T}$ for $\gamma\gamma \to qV\bar{q}\bar{V}$. Their general structure is

$$\hat{T}_{\{\lambda\}}(\hat{t},\hat{u}) = C_{\mathrm{cf}}^{(3)}\,\overline{T}_i^{(3,V)}(\hat{t},\hat{u})F_V^{(3)} + C_{\mathrm{cf}}^{(4)}\,\overline{T}_i^{(4,D)}(\hat{t},\hat{u})F_V^{(4)} + C_{\mathrm{cf}}^{(5)}\,\overline{T}_i^{(5,V)}(\hat{t},\hat{u})F_V^{(5)}, \quad (13)$$

where $C_{\mathrm{cf}}^{(n)}$ are the appropriate charge-flavor factors. The subscript $i = 1,\dots,13$ labels the helicity-combinations according to Eq. (2). Each n-point contribution $\overline{T}^{(n,V)}$

is found from a separately gauge-invariant set of Feynman diagrams, where $(n-2)$ gauge bosons couple to the diquark. The $\overline{T}^{(n,V)}$ are multiplied with the appropriate diquark form factors $F_V^{(n)}$, parameterizing the composite nature of diquarks. For further details we refer to [5].

The analytical results for $\overline{T}_i^{(n,V)}$ are presented in the following. For their calculation we employed the algebraic computer program *Mathematica* [24] with the package `FeynCalc` [25]. We do not list n-point contributions that are suppressed by at least $\mathcal{O}(m_{\mathrm{B}_{10}}^2/\hat{s})$, since they are neglected in our numerical calculations. Those include, for example, all 5-point functions $\overline{T}_i^{(5,V)}$ and all amplitudes with (anti)baryon helicity $\pm\frac{3}{2}$. Note that the parameterization of the form factors $F_V^{(n)}$ for 4- and 5-point functions provides additional inverse powers of $\hat{s}$ as compared to $F_V^{(3)}$. We abbreviate $C=(4\pi)^2 C_F \alpha\,\alpha_s$, where $C_F=\frac{4}{3}$ is the color factor, and α denotes the fine structure constant $\alpha \approx 1/137$. κ_{V} is the anomalous magnetic moment of the vector diquark.

$$
\begin{aligned}
\overline{T}_1^{(3,V)}(\hat{t},\hat{u}) &= -\frac{4}{3}C\frac{\kappa_{\mathrm{V}}}{m_{\mathrm{B}_{10}}^2\sqrt{\hat{u}\hat{t}}}\left(\frac{\hat{u}}{x_1y_1}+\frac{\hat{t}}{x_2y_2}\right)\\
\overline{T}_2^{(3,V)}(\hat{t},\hat{u}) &= \frac{2}{3}C\frac{1}{m_{\mathrm{B}_{10}}}\frac{\sqrt{\hat{s}}\,\hat{s}}{\hat{u}\hat{t}}\frac{x_1+y_1}{x_1y_1}\\
\overline{T}_4^{(3,V)}(\hat{t},\hat{u}) &= C\frac{2}{3\,m_{\mathrm{B}_{10}}}\frac{1}{\sqrt{\hat{s}}\,\hat{u}\hat{t}}\frac{1}{x_1x_2y_1y_2}\Bigg\{\\
&\quad -\kappa_{\mathrm{V}}\left[(2x_1-3)y_2\hat{t}^2+(2y_1-3)x_2\hat{u}^2-4x_1y_1\hat{u}\hat{t}\right]+\\
&\quad +\left[(x_2+y_2)\left(\hat{t}^2x_1y_2+\hat{u}^2x_2y_1-2x_1y_2\hat{u}\hat{t}\right)-y_2\hat{t}^2-x_2\hat{u}^2\right]\Bigg\}\\
\overline{T}_5^{(3,V)}(\hat{t},\hat{u}) &= \overline{T}_1^{(3,V)}(\hat{u},\hat{t})\\
\overline{T}_6^{(3,V)}(\hat{t},\hat{u}) &= -C\frac{2}{3}\frac{1}{m_{\mathrm{B}_{10}}}\frac{\sqrt{\hat{s}}\,\hat{s}}{\hat{u}\hat{t}}\left[-(1+\kappa_{\mathrm{V}})\frac{x_1y_2^2+y_1x_2^2}{x_1x_2y_1y_2}+\kappa_{\mathrm{V}}\frac{x_1+y_1}{x_2y_2}\right]\\
\overline{T}_2^{(4,V)}(\hat{t},\hat{u}) &= -\frac{2}{3}C\frac{\kappa_{\mathrm{V}}(1-\kappa_{\mathrm{V}})\sqrt{\hat{s}}}{m_{\mathrm{B}_{10}}^3}\frac{1}{x_1x_2^2y_1y_2^2}.
\end{aligned}
$$

The hard-scattering amplitudes for Compton scattering off decuplet baryons are related to the amplitudes listed above via crossing [22]. The corresponding elementary helicity amplitudes, $\gamma q D \to \gamma q D$ have been computed separately as a check. They can be obtained from the authors upon request.

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References

[1] M. R. Pennington, Nucl. Phys. Proc. Suppl. **82**, 291 (2000) [hep-ph/9907353]; S. J. Brodsky, eConf **C0309101**, WEPL001 (2003) [hep-ph/0311355].

[2] S. J. Brodsky, *High energy photon photon collisions at a linear collider*, hep-ph/0404186.

[3] P. Achard *et al.* [L3 Collaboration], Phys. Lett. B **536**, 24 (2002) [hep-ex/0204025];
G. Abbiendi *et al.* [OPAL Collaboration], Eur. Phys. J. C **28**, 45 (2003) [hep-ex/0209052];
C. C. Kuo, S. C. Kao, A. Chen, W. T. Chen and S. Uehara [Belle Collaboration], Nucl. Phys. A **721**, 817 (2003);
P. Achard *et al.* [L3 Collaboration], Phys. Lett. B **571**, 11 (2003) [hep-ex/0306017];
T. Barillari, Nucl. Phys. Proc. Suppl. **126**, 301 (2004) [hep-ex/0307066];
D. Urner, AIP Conf. Proc. **698**, 566 (2004) [hep-ex/0309045].

[4] M. Diehl, P. Kroll and C. Vogt, Eur. Phys. J. C **26**, 567 (2003) [hep-ph/0206288]; P. Kroll, Fizika B **13**, 153 (2004) [hep-ph/0310327].

[5] C. F. Berger and W. Schweiger, Eur. Phys. J. C **28**, 249 (2003) [hep-ph/0212066].

[6] M. Karliner, in *Proc. of the e^+e^- Physics at Intermediate Energies Conference* ed. Diego Bettoni, eConf **C010430**, W10 (2001) [hep-ph/0108106].

[7] G. R. Farrar, E. Maina and F. Neri, Nucl. Phys. B **259**, 702 (1985) [Erratum-ibid. B **263**, 746 (1986)].

[8] M. Anselmino, F. Caruso, P. Kroll and W. Schweiger, Int. J. Mod. Phys. A **4**, 5213 (1989).

[9] G. R. Farrar, H. Zhang, A. A. Ogloblin and I. R. Zhitnitsky, Nucl. Phys. B **311**, 585 (1989).

[10] M. Karliner and S. Nussinov, Phys. Lett. B **538**, 321 (2002) [hep-ph/0202234].

[11] G. P. Lepage and S. J. Brodsky, Phys. Rev. D **22**, 2157 (1980);
S. J. Brodsky and G. P. Lepage, Adv. Ser. Direct. High Energy Phys. **5**, 93 (1989).

[12] A. V. Efremov and A. V. Radyushkin, Phys. Lett. B **94**, 245 (1980).

[13] M. Anselmino, P. Kroll and B. Pire, Z. Phys. C **36**, 89 (1987).

[14] R. Jakob, P. Kroll, M. Schürmann and W. Schweiger, Z. Phys. A **347**, 109 (1993) [hep-ph/9310227].

[15] P. Kroll, M. Schürmann and W. Schweiger, Int. J. Mod. Phys. A **6**, 4107 (1991);
P. Kroll, T. Pilsner, M. Schürmann and W. Schweiger, Phys. Lett. B **316**, 546 (1993) [hep-ph/9305251].

[16] C. F. Berger, B. Lechner and W. Schweiger, Fizika B **8**, 371 (1999) [hep-ph/9901338];
C. F. Berger, B. Jäger and W. Schweiger, Nucl. Phys. A **689**, 390 (2001) [hep-ph/0009295];
B. Jäger and W. Schweiger, Nucl. Phys. A **711**, 203 (2002) [hep-ph/0206189];
C. F. Berger and W. Schweiger, Fizika B **13**, 74 (2004) [hep-ph/0309095].

[17] C. F. Berger and W. Schweiger, Phys. Rev. D **61**, 114026 (2000) [arXiv:hep-ph/9910509].

[18] H. Albrecht *et al.* [ARGUS collaboration], Z. Phys. C **42**, 543 (1989).

[19] M. Anselmino, F. Murgia and F. Caruso, Phys. Rev. D **42**, 3218 (1990).

[20] W. Rarita and J. S. Schwinger, Phys. Rev. **60**, 61 (1941).

[21] V. L. Chernyak, A. A. Ogloblin and I. R. Zhitnitsky, Z. Phys. C **42**, 569 (1989) [Yad. Fiz. **48**, 1410 (1988 SJNCA,48,896-904.1988)].

[22] C. Bourrely, J. Soffer and E. Leader, Phys. Rept. **59**, 95 (1980);
E. Leader, Cambridge Monogr. Part. Phys. Nucl. Phys. Cosmol. **15**, 1 (2001).

[23] B. Echenard, talk at *Photon 2003*, unpublished, and private communication.

[24] S. Wolfram, *The Mathematica Book*, 5th Edition, Wolfram Media, 2003.

[25] R. Mertig, M. Böhm and A. Denner, Comput. Phys. Commun. **64**, 345 (1991).