BLEJSKE DELAVNICE IZ FIZIKE LETNIK 19, ŠT. 1
BLED WORKSHOPS IN PHYSICS VOL. 19, NO. 1

ISSN 1580-4992

**Proceedings of the Mini-Workshop**

# Double-charm baryons and dimesons

**Bled, Slovenia, June 17 – 23, 2018**

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DMFA – ZALOŽNIŠTVO
LJUBLJANA, NOVEMBER 2018

## The Mini-Workshop *Double-charm baryons and dimesons*

**was organized by**

*Society of Mathematicians, Physicists and Astronomers of Slovenia*
*Department of Physics, Faculty of Mathematics and Physics, University of Ljubljana*

**and sponsored by**

*Department of Physics, Faculty of Mathematics and Physics, University of Ljubljana*
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# Preface

This year the Bled Mini-Workshop was devoted to the double-charm hadrons, as the discovery of the double-charm baryon $\Xi_{cc}^{++} = ccu$ at LHCb revived the interest for the search of double charm dimesons. The first question is whether the $DD^*$ dimeson is bound or is it a low-lying resonance. In models using quark exchange it is only weakly bound; pion exchange might supply extra binding but no reliable model calculation has been performed so far. The second question is what would be a good signature for its detection, since the production and decay processes are not as clear as in the case of $\Xi_{cc}^{++}$. We hope that our present discussions and future ideas could give a hint to experimentalists for a search, especially at the upgraded Belle II facility.

Both constituent quark models and lattice QCD simulations are helpful and are already giving a lot of insight. They both support the value 3620 MeV for the mass of the $\Xi_{cc}^{++}$ baryon found at LHCb. However, the controversial $\Xi_{cc}^{+}$ baryon at 3520 MeV (SELEX) is not yet completely excluded as a possible lower state. The mechanism of the dimeson production via a cc diquark is a challenge. The encouragement comes from the copious double $c\bar{c}$ production seen at Belle. Theoretical explanation is still missing.

Nucleon resonances were also a side topic of the meeting, among them mostly the Roper resonance which remains a major challenge in spite of recent developments in lattice QCD. It is broad and hard to see directly in spectra, and it is still unclear to what extent it is a breathing mode of the proton (a three-quark system) or a dynamically generated resonance. Experimentally it has been studied by electroexcitation, and theoretically by coupled-channel analysis (the $\pi N$, $\sigma N$, $\pi\Delta$ channels, as well as channels with $\rho, \omega, \eta$ and K). The provocative peaks E(38 MeV) and Z(57.5 GeV) also ask for explanation.

We would like to thank again all participants for coming and providing new perspectives on the phenomena of our common interest. It is so encouraging that in such a small and friendly group we can tell each other the strong and the weak points of our approaches and profit from the frank criticism and suggestions. We shall do our best that these popular Mini-Workshops continue every year and we hope to see you at Bled again.

Ljubljana, November 2018 B. Golli, M. Rosina, S. Širca

P. S.: Several figures and diagrams originally presented in colour are reproduced here in grayscale, whereby some information is lost. The color version is available at `http://www-f1.ijs.si/BledPub`.

# Predgovor

Letos je bila blejska delavnica posvečena dvojno čarobnim hadronom, saj je odkritje dvojno čarobnega bariona $\Xi_{cc}^{++} = ccu$ na pospeševalniku LHCb oživilo zanimanje za iskanje dvojno čarobnih dimezonov. Prvo vprašanje je, ali je dimezon $DD^*$ vezan, ali je nizko ležeča resonanca. Če upoštevamo izmenjavo kvarkov, je samo šibko vezan; izmenjava pionov lahko dodatno veže, toda doslej še ni zanesljivega računa. Drugo vprašanje je, kako bi ga prepoznali, kajti tvorba in razpad nista tako značilna kot pri barionu $\Xi_{cc}^{++}$. Upamo, da bodo naša sedanja razmišljanja in bodoče ideje pomagale eksperimentalcem pri iskanju, zlasti na povečanem detektorju Belle II.

Dosti vpogleda dobimo od računov s kvarkovimi modeli kakor tudi od simulacije kromodinamike na mreži. Oboje je v skladu z maso bariona $\Xi_{cc}^{++}$ 3620 MeV, ki so jo izmerili na pospeševalniku LHCb. Toda sporna masa bariona $\Xi_{cc}^{+}$ okrog 3520 MeV pri detektorju SELEX še ni čisto izključena kot možno najnižje stanje. Izziv je tudi produkcija dimezonov preko dikvarka cc. Vzpodbudo daje obilna produkcija dvojnih parov $c\bar{c}$, ki so jo izmerili na detektorju Belle. Teoretična razlaga še manjka.

Kot stransko temo srečanja smo obravnavali tudi nukleonske resonance, zlasti Roperjevo resonanco, ki predstavlja še vedno velik izziv kljub nedavnemu napredku pri kromodinamiki na mreži. Resonanca je široka in jo je težko opaziti v v spektrih; še vedno je nejasno, do kolikšne mere je sistem treh kvarkov ("dihanje" protona) ali dinamično povzročena resonanca. Eksperimentalno so jo preučevali z vzbujanjem z elektroni, teoretično pa z analizo sklopljenih kanalov ( $\pi N$, $\sigma N$, $\pi\Delta$), tudi takih z mezoni $\rho, \omega, \eta$ in K. Izzivalna vrhova pri 38 MeV in pri 57,5 GeV tudi kličeta k preverjanju in razlagi.

Radi bi se zahvalili vsem udeležencem, da so se udeležili srečanja in nudili nove perspektive pri pojavih, ki nas vse zanimajo. Vzpodbudno je, da si lahko v takšni majhni in prijateljski skupini povemo tako močne kot šibke točke pri naših pristopih in nam koristijo odkrita kritika in namigi. Potrudili se bomo, da se bo ta priljubljena Delavnica nadaljevala vsako leto in upamo, da se spet vidimo na Bledu.

Ljubljana, november 2018 B. Golli, M. Rosina, S. Širca

P. S. Marsikatere slike in diagrame smo prejeli v barvah, toda v tiskanem Zborniku so sivi, s čimer se zgubi nekaj informacije. Barvno verzijo lahko dobite na `http://www-f1.ijs.si/BledPub`.

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# Mesonic Spectra: Experimental Data and Their Interpretation*

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**Abstract.** We discuss experimental data and their interpretation. In particular, we argue that spectra of quark-antiquark systems should better be studied from configurations with well-defined quantum numbers, the most suitable system being charmonium. We suggest probable future findings based on the existing low-statistics data for charmonium and bottomonium. We also briefly review our findings for the E(38 MeV) and Z(57.5 GeV) bosons.

Observation and its interpretation are human activities not restricted to a handful of experts but open to anyone who feels the need to express an opinion. A nice example of the importance of observation is the meticulous registration of atomic and molecular line spectra during the nineteenth and the twentieth century. Moreover, its history of interpretation shows exemplarily the struggle of the human mind to escape from prevailing standard models. It starts with the color-separation theory of Wollaston, based on his pioneering observation in the early 1800s of seven dark lines in the solar spectrum. Decades later that interpretation was proven to be wrong by Kirchhoff and Bunsen, based on the observation of emission spectra. Thomson's plum-pudding model in the early 1900s, shortly after the discovery of electrons, was the last attempt to keep observation within accepted theories. Finally, Bohr's proposal gave the breakthrough for a solid description of line spectra and the emergence of a new standard model. A century full of observation, improving equipment and new discoveries had passed in order to figure it out.

The history of atomic and molecular line spectra resembles that of mesonic spectra. But it fails when it comes to high-quality data. Bohr's model could be tested on a wealth of experimental results. Models for mesonic resonances do not have such luxury at their disposal, which has culminated in a plethora of

* Talk presented by E. van Beveren

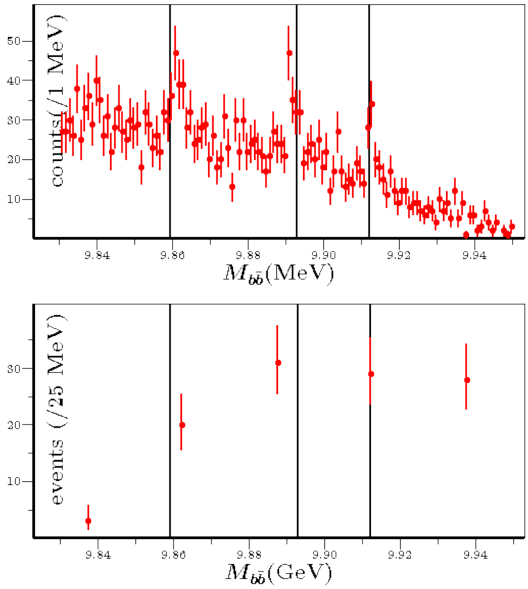


**Fig. 1.** The 1P beautonium states. Top: ARGUS data (DESY, 1985). Bottom: ATLAS data (CERN, 2011).

speculations. Though one wishes that future experiments will improve on statistics, the reality is quite different, as is most strikingly exhibited in Fig. 1, where three-decades-old data [1] on $b\bar{b} \to \Upsilon\gamma$ are compared to more recent results [2]. In this short paper we will highlight some of our somewhat speculative suggestions about the interpretation of mesonic spectra based on observation but not yet confirmed by dedicated experiments.

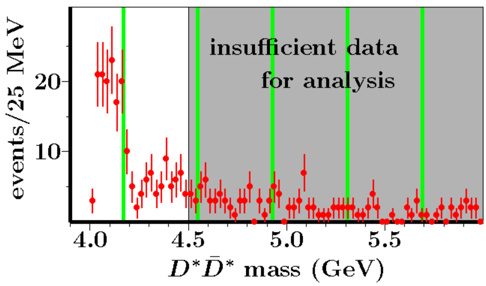


**Fig. 2.** The $D^*\bar{D}^*$ mass distribution measured and published by the BABAR Collaboration.

At several occasions we have pointed out the indispensable need for high-statistics data on two-particle mass distributions. As an example may serve the data shown in Fig. 2, where we represent a $D^*\bar{D}^*$ mass distribution measured and published by the BABAR Collaboration [3]. At first sight these data do not give us further information on the $c\bar{c}$ vector-meson spectrum. Indeed, a bin size of 25 MeV is clearly too large for the narrow dominantly-D states and even for the somewhat broader dominantly-S states, whereas also the number of events is barely enough to show an enhancement of the $\psi(4040)$. However, one must bear in mind the following.

In the first place, the reconstruction of a pair of $D^*$ mesons out of kaons, pions and photons is a far from trivial task. The procedure is indicated in Ref. [3]. But it is not clear to us what fraction of produced $D^*\bar{D}^*$ pairs is recognized that way. We assume that it is a relatively small fraction.

Next, we know from theory that the higher the $c\bar{c}$ vector meson mass, the smaller its coupling to $D^*\bar{D}^*$. The reason is that, under the assumption of ${}^3P_0$ quark-pair creation, the number of possible two-meson configurations to which a $c\bar{c}$ vector meson couples grows rapidly with radial excitation [4]. Consequently, the coupling to a specific channel, in the present case $D^*\bar{D}^*$, diminishes substantially for higher radial excitations, thus leading to decreasing enhancements.

Finally, S- and D-wave $c\bar{c}$ vector states mix, which implies that pure S- or D-wave states do not exist in nature. But mixing also has two other interesting consequences. Namely, the dominantly D-wave states almost decouple from meson-pair production, leading to narrow resonances and small mass shifts, whereas the relatively broad S-wave states can easily dominate in decay and so confound their classification. The second consequence of mixing is that the dominantly S-wave states couple more strongly to meson.pair production than expected for pure S-wave states, giving rise to larger widths and considerable mass shifts [5].

So the question comes up why, in the absence of good data, we insist on dealing with $c\bar{c}$ vector mesons. The answer to that question rests in our belief that these mesons form the backbone of quark-antiquark $q\bar{q}$ spectra:

1. In the process $e^-e^+ \to D^*\bar{D}^*$, vector-meson dominance ensures the production of $c\bar{c}$ vector states. Hence, there is no confusion with different quantum numbers.

2. Little to no influence is expected from non-strange, strange and bottom $q\bar{q}$ pairs.

Consequently, when we know the full details of the $c\bar{c}$ vector spectrum, we can easily fill up the gaps for the remaining configurations and then use that for the analyses of different flavor combinations.

In Fig. 3 we have depicted the poor data for the $D^*\bar{D}^*$ mass distribution, together with a comparison to our predictions [6]. The crosses on the horizontal axis indicate the masses of bare $c\bar{c}$ vector states, *i.e.*, the spectrum in the absence of two-meson configurations, where in our model [5] S- and D-states are degenerate. By allowing $c\bar{c}$ to couple to open-charm configurations, the predicted dominantly-D states shift only a few MeV, whereas the mass shifts for the dominantly-S states are of the order of 100-300 MeV. The enhancement indicated by $\Lambda_c\Lambda_c$ is explained in Ref. [7] (see also Fig. 4). Given the importance of the $c\bar{c}$ vector states for meson spectroscopy, it escapes us why after four decades high-statistics data still do not exist. But maybe Fig. 5 explains it. In the following we will make some suggestions about the $b\bar{b}$ vector spectrum, as well as the E(38 MeV) and Z(57.5 GeV) bosons.

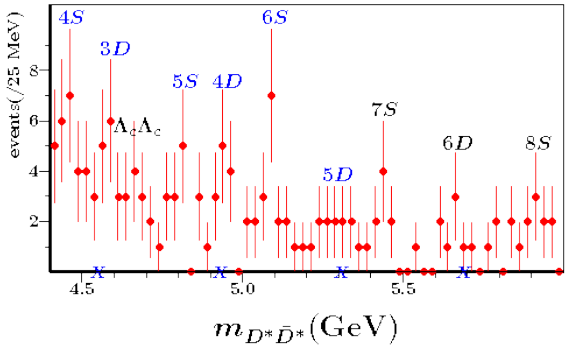


**Fig. 3.** The poor data for the $D^*\bar{D}^*$ mass distribution together with a comparison to our predictions.

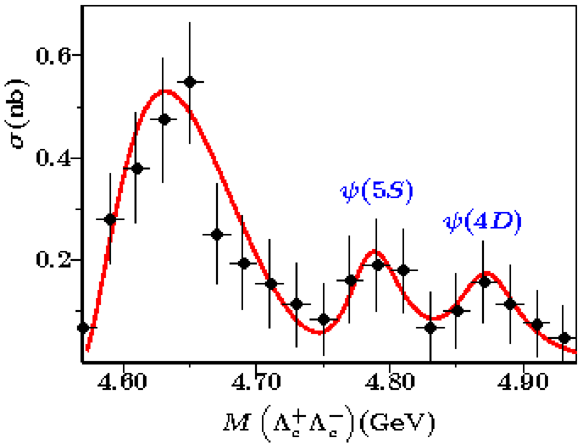


**Fig. 4.** $\psi(5S)$ and $\psi(4D)$ besides the large signal ($Y(4660)$) at the $\Lambda_c\Lambda_c$ threshold.

In Fig. 6 we show our result for the $\Upsilon(2D)$ $b\bar{b}$ vector state at about 10.5 GeV, some 70 MeV below the $B\bar{B}$ threshold. The data are taken from Ref. [8], while our analysis is discussed in Ref. [9]. A bound state as close to threshold as the $\Upsilon(2D)$ is supposed to have a large influence on the threshold enhancement. In Fig. 7 we have depicted $R_b$-ratio data from Ref. [10], in which one indeed observes a large threshold enhancement peaking at about 10.58 GeV, followed by two more modest enhancements above the $B\bar{B}^*$ and $B^*\bar{B}^*$ thresholds. The figure also shows that the former enhancement is listed under $\Upsilon(4S)$ in Ref. [11] and, moreover, that our model does not agree with that assignment. This is substantiated in Fig. 8, where TE stands for threshold enhancement, BW for Breit-Wigner line shape. In view of the above discussion on the decrease of resonance enhancements, it seems to us quite reasonable that the $\Upsilon(4S)$ is a modestly peaked structure. Moreover, its central mass at about 10.73 GeV agrees better with our model predictions.

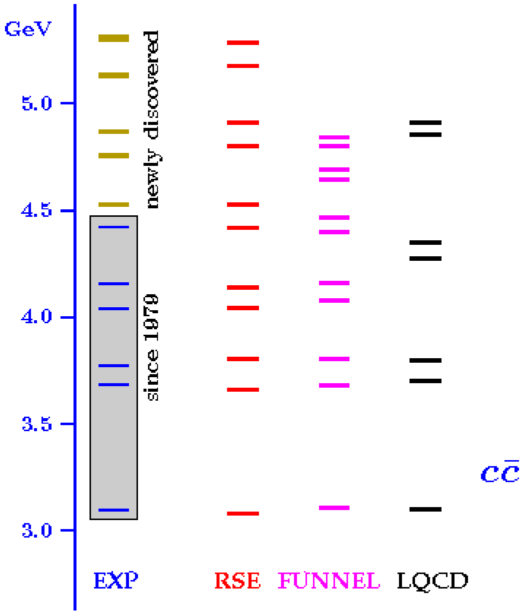


**Fig. 5.** The 2009 $c\bar{c}$ vector spectrum. EXP: PDG + new states; RSE: PRD **21**, (1980); FUNNEL: representative for MOST other models; LQCD: representative for lattice QCD.

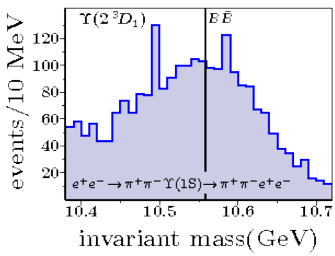


**Fig. 6.** Our result for the $\Upsilon$(2D) $b\bar{b}$ vector state at about 10.5 GeV, some 70 MeV below the $B\bar{B}$ threshold.

The discovery of the E(38) boson [12] is discussed in the web version of the talk [13] (click `start`, then `E38`, and check the slides from `r0` to `compass2`). The slides `r0` and `ρ0` show why we expected a quantum of about 30–40 MeV, to be associated with quark-pair creation, already since the 1980s. Hints from experimental results came later as exhibited in slides from `wobbles` to `more`. More promising data [14, 15] are shown in slides from $\gamma\gamma$ to `compass2`. However, the COMPASS Collaboration contested our proposal by claiming that the enhancement at about 38 MeV is due to an artifact, the details of which are explained in Ref. [16]. Now it must be mentioned that the COMPASS Collaboration has done excellent work on light-meson spectroscopy [17]. Unfortunately, in the effort to substantiate the artifact claim, the COMPASS Collaboration compared apples and

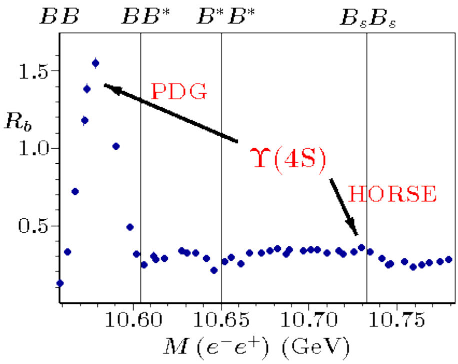


**Fig. 7.** The $R_b$-ratio data from Ref. [10] with a large threshold enhancement peaking at about 10.58 GeV, followed by two more modest enhancements above the $B\bar{B}^*$ and $B^*\bar{B}^*$ thresholds.

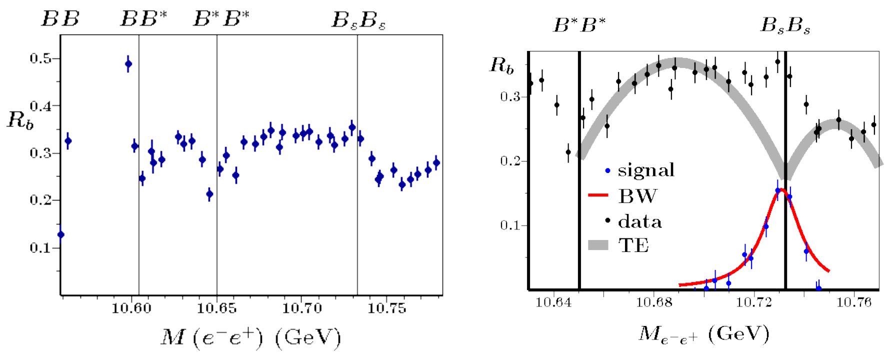


**Fig. 8.** Left: detail enhancements. Right: threshold enhancements and resonances.

oranges by referring to the low-statistics data of Ref. [14] instead of to the high-statistics data of Ref. [15]. But even the Monte-Carlo simulation for the former data do not minimally confirm their explanation, as shown in Fig. 9. We are still awaiting the follow-up on Ref. [18], which tentatively confirmed the existence of the E(38) [19].

For our suggestion of the existence of a boson at about 57.5 GeV, we only see some hints in experimental data. In Fig. 10 [20] we observe a rather sharp dip in the amplitude at about 115 GeV. When we sift through other observations in that energy region from the CMS, ATLAS and LEP Collaborations and combine the data [20–23] in Fig. 14, we find some indications of agreement. Now, such a sharp minimum in the data could indicate the onset of a threshold enhancement, moreover inflated due to the presence of a resonance at about 125 GeV, and most probably resulting from the creation of a pseudoscalar (or scalar) boson pair of half the onset mass each The L3 Collaboration might have searched for such a boson in $Z \to \gamma\gamma\gamma$ [24]. But with a total of 87 events not much statistics can be expected, as we see in Figs. 11 and 12. Nevertheless, a small effect is visible

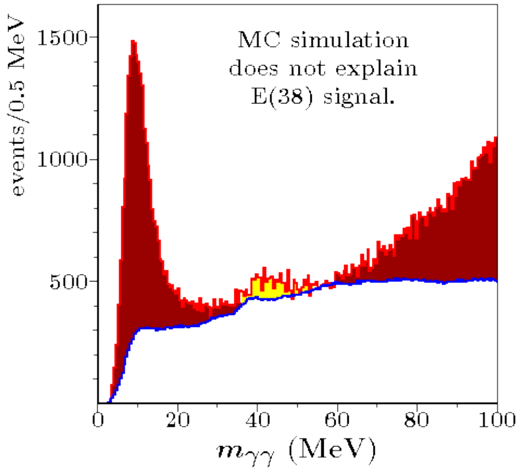


**Fig. 9.** Diphoton Monte-Carlo simulation from the COMPASS Collaboration compared to the data.

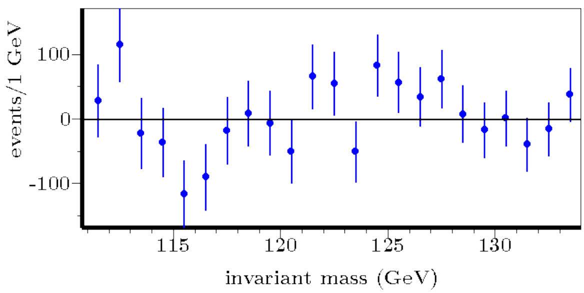


**Fig. 10.** Diphoton data from the CMS Collaboration.

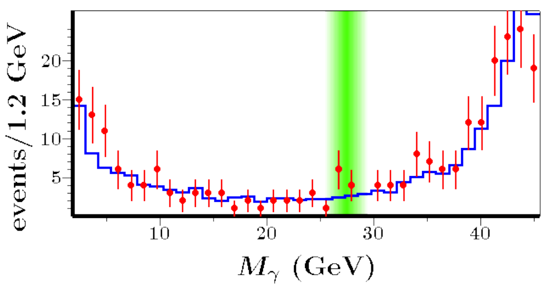


**Fig. 11.** $Z \to \bar{Z}\gamma \to \gamma\gamma\gamma$. Solid line: QED expectation. The shaded area represents the expected one-photon energy for $M(\bar{Z}) = 57.5$ GeV.

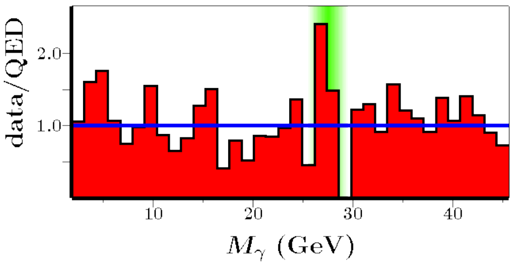


**Fig. 12.** Data divided by QED. The excess is where expected for $\bar{Z}$ of 57.5 GeV.

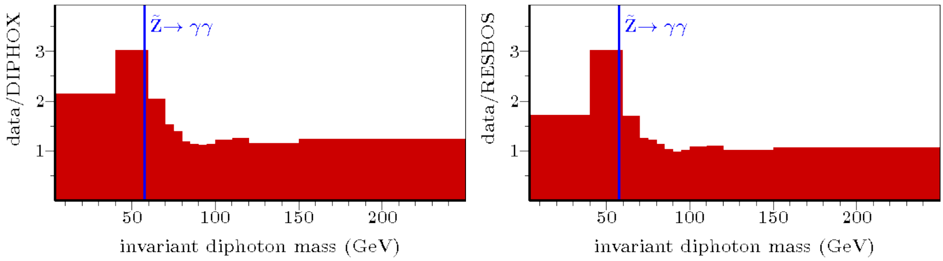


**Fig. 13.** Diphoton data from the CMS Collaboration compared to Standard-Model predictions DIPHOX (left) and RESBOS (right).

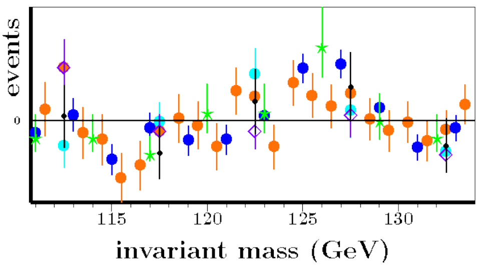


**Fig. 14.** Other LHC and LEP data agree. Shown are the data for CMS $\gamma\gamma$, ATLAS $\gamma\gamma$, CMS 4 leptons, ATLAS 4 leptons, L3 $e^+e^- \to \tau^+\tau^-(\gamma)$ and L3 $e^+e^- \to \mu^+\mu^-(\gamma)$.

precisely where expected. Further indications come from comparison of diphoton data from the CMS Collaboration with predictions of DIPHOX and RESBOS [25], shown in Fig. 13. It should not be too difficult to obtain clean data at LHC to improve the $Z \to 3\gamma$ statistics.

## Acknowledgment

One of us (EvB) wishes to thank the organizers of the Mini-Workshop Bled 2018 on *"Double-charm baryons and dimesons"* for their hospitality.

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VOL. 19, NO. 1
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*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Heavy baryon spectroscopy from lattice QCD

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**Abstract.** In this report, the most recent and precise estimates of masses of ground state baryons using lattice QCD are discussed. Considering the prospects in the heavy baryon sector, lattice estimates for these are emphasized. The first and only existing lattice determination of the highly excited $\Omega_c$ excitations in relation to the recent LHCb discovery is also discussed.

## 1 Introduction

Since its inception, heavy hadron physics continues to be in the limelight of scientific interests in understanding the nature of strong interactions. While heavy mesons have been studied extensively both experimentally and theoretically [1–3], studies on heavy baryons remained dormant. In this respect, the year 2017 featured two important landmarks in the heavy baryon physics. First of this is the unambiguous observation by LHCb collaboration of five new narrow $\Omega_c$ resonances in $\Xi_c^+ K^-$ invariant mass distribution in the energy range between $3000 - 3120$ MeV [4]. Four out of these five resonances were later confirmed by Belle collaboration [5]. Second landmark is the discovery of a doubly charmed baryon, $\Xi_{cc}^+(ccu)$ with a mass of $3621.40 \pm 0.78$ MeV by LHCb Collaboration [6]. Anticipating the discovery of many more hadrons (including baryons) from the huge data being collected at LHCb and Belle II, heavy hadron spectroscopy using *ab-initio* first principles methodology such as lattice QCD is of great importance.

Lattice QCD has been proven to be a novel non-perturbative technology in investigating the physics of low energy regime of QCD. Remarkable progress has been achieved over past ten years in making large volume simulations with physical quark masses, impressive statistical precision and good control over the systematic uncertainties [7–10]. In this report, a collection of lattice determinations of baryon masses that are well below allowed strong decay thresholds are summarized. A recent and only existing calculation of excited $\Omega_c$ baryons is discussed and a qualitative comparison with the experiment is made.

## 2 Lattice methodology

Hadron spectroscopy on the lattice proceeds through evaluation of Euclidean two point correlation functions,

$$C_{ij}(t_f - t_i) = \langle O_i(t_f) O_j^\dagger(t_i) \rangle = \sum_{n=1} \frac{Z_i^n Z_j^{n*}}{2E_n} e^{-E_n(t_f - t_i)}, \tag{1}$$

between different hadronic currents ($O_i(t)$) that are carefully built to respect the quantum numbers of interest. A generic baryon current or interpolator has a structure

$$O_i(\mathbf{x}, t) = \epsilon_{abc} S_i^{\alpha\beta\delta}(\mathbf{x}) q^a_{1,\alpha}(\mathbf{x}) q^b_{2,\beta}(\mathbf{x}) q^c_{3,\delta}(\mathbf{x}), \tag{2}$$

where $q_j$ are the quark fields, $\epsilon$ is the color space anti-symmetrizing Levi-Civita tensor and S carries all the flavor and spatial structure of the interpolator that determines the quantum information. $C_{ij}(t_f - t_i)$ are evaluated on lattice QCD ensembles that are generated via Monte Carlo techniques. A general practice is to compute matrices of correlation functions between a basis of carefully constructed interpolating currents $O_i(t)$ and solving the generalized eigenvalue problem (GEVP) [11–13]

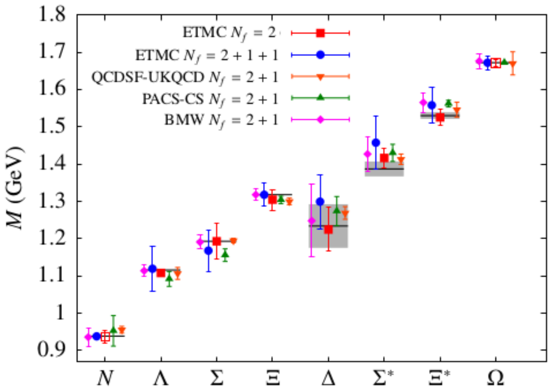
$$C_{ij}(t) v_j^n(t - t_0) = \lambda^n(t - t_0) C_{ij}(t_0) v_j^n(t - t_0). \tag{3}$$

Hadron energies ($E_n$) are extracted from non-linear fits to the large time behavior of the eigenvalues $\lambda^n(t - t_0)$. The eigenvectors ($v_j^n(t - t_0)$) are related to the operator state overlaps ($Z_i^n = \langle O_i | n \rangle$) that carry the quantum information of the propagating state. Basic principles remain the same as above, while details of the methodology differ between different groups in the lattice community. e.g. lattice ensembles being used in the study, lattice formulation of action for the fermion and the gauge fields, the hadron interpolators, different degree of control over the lattice systematics, etc. The success of lattice investigations are reflected in mutual agreement of the results they provide and their agreement with experiments.

All results presented in this report are estimated within the single hadron approximation, where only three quark interpolators (as in eqn. 2) are considered in the analysis and neglects effects of any nearby strong decay thresholds. This is a justifying assumption for most of the baryons discussed in this report, considering the fact that all of them are deeply below the respective lowest strong decay thresholds. Results for those baryons, which might be influenced by any nearby threshold effects will be alerted in the respective discussions.

## 3 Results

*Light, strange and singly charm baryons* : We begin our discussion with some benchmark calculations of baryon ground states that are experimentally well determined. In Fig. 1, a summary of lattice QCD estimates for the positive parity light baryon ground states (figure adapted from Ref. [10]) are presented at the top and for positive parity singly charm baryon ground states are shown at the bottom. Most of the baryons being discussed are deeply bound and stable to strong decays. Their masses as determined from the discrete energy spectrum on the lattice agree quite well with experiments. Agreement between all the lattice estimates with varying degree of control over the systematics involved in respective calculations and with the experiments demonstrate the power of lattice QCD techniques in making reliable predictions. However, lattice estimates for masses of baryon resonances, such as $\Delta$, $\Sigma^*$ and $\Xi^*$ that can decay strongly, are less rigorous. They demand a computation of correlation matrices build out of baryon



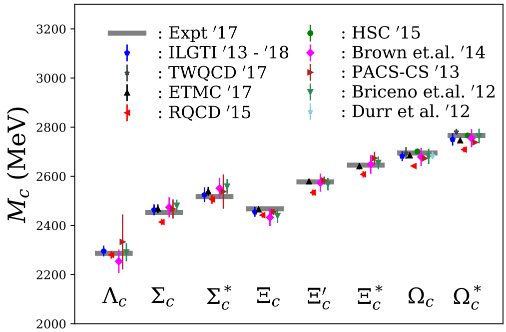




**Fig. 1.** Top: summary of lattice estimates for positive parity light and strange baryons from selected lattice investigations - ETMC $N_f$=2 [10], ETMC $N_f$=2+1+1 [14], QCDSF-UKQCD $N_f$=2+1 [15], PACS-CS $N_f$=2+1 [8] and BMW $N_f$=2+1 [7]. Bottom: summary of lattice estimates for positive parity singly charm baryons : ILGTI '13-'18 [16, 17], TWQCD '17 [18], ETMC '17 [10], RQCD '15 [19], HSC '15 [20, 21], Brown *et al* '14 [22], PACS-CS '13 [8], Briceño *et al* '12 [23], Dürr *et al* '12 [24].

interpolators (as in eqn. 2) plus baryon-meson interpolators (corresponding to the allowed strong decay modes). The masses of baryon resonances then have to be inferred from the infinite volume scattering matrices build from the discrete spectrum extracted from such correlation matrices. Such investigations are being practised extensively by many collaborations to understand various mesonic resonances (see Ref. [3]), while existing lattice investigations of baryon resonances in this direction are limited to a few [25, 26].

*Doubly heavy baryons* : In Fig. 2, a summary of lattice QCD estimates for positive parity doubly charm baryon ground states at the top is presented. For the $\Xi_{cc}(1/2^+)$ baryon, good agreement between all lattice estimates (all of which predates the LHCb-discovery [6]) and with LHCb estimate is quite evident from the figure. At this point, the reader is reminded of the observation of another baryon resonance by SELEX collaboration in 2002 [27] at a mass of 3519(1) MeV,

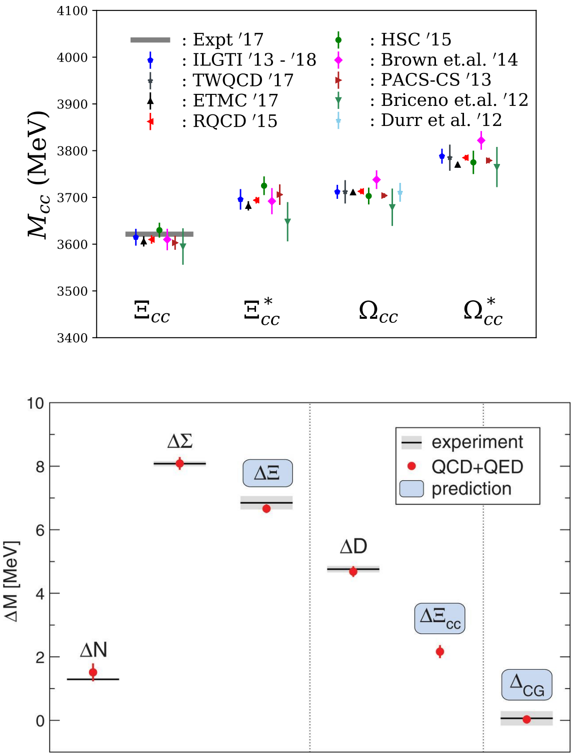



**Fig. 2.** Top: summary of lattice estimates for positive parity doubly charm baryons. References as given in Fig. 1 caption. Bottom: Hadron isospin splittings as determined by BMW collaboration [9].

which is addressed as a $\Xi_{cc}(1/2^+)$ baryon. All lattice estimates, being well above this energy, disfavors this observation. The bottom figure shows a summary of baryon isospin splittings as calculated by BMW collaboration [9]. This calculation involved lattice QCD and QED computations with four non-degenerate fermion flavors to estimate the isospin mass splitting in the nucleon, Σ, Ξ, D and $\Xi_{cc}$ isospin multiplets. Precise estimation of the neutron-proton isospin splitting and the other known splittings demonstrate the reliability of these estimates. In this calculation, the isospin splitting of $\Xi_{cc}(1/2^+)$ baryon was estimated to be 2.16(11)(17) MeV. This excludes the possibility that LHCb and SELEX candidates for $\Xi_{cc}(1/2^+)$ baryon are isospin partners.

Estimates for other doubly charm baryons, that are yet to be discovered, can also be observed to be very well determined and consistent between different lattice calculations from the top of Fig. 2. Anticipating a near future discovery of the charmed-bottom hadrons at LHCb, at the top of Fig. 3 lattice predictions for such hadrons from a recent investigation [28] are shown. The lattice prediction for only know charmed-bottom hadron, $B_c$ meson, is found to be in agreement with the experiment, while the lattice predictions for other channels considered are consistent with another preceding calculation [22] with less control over systematics.

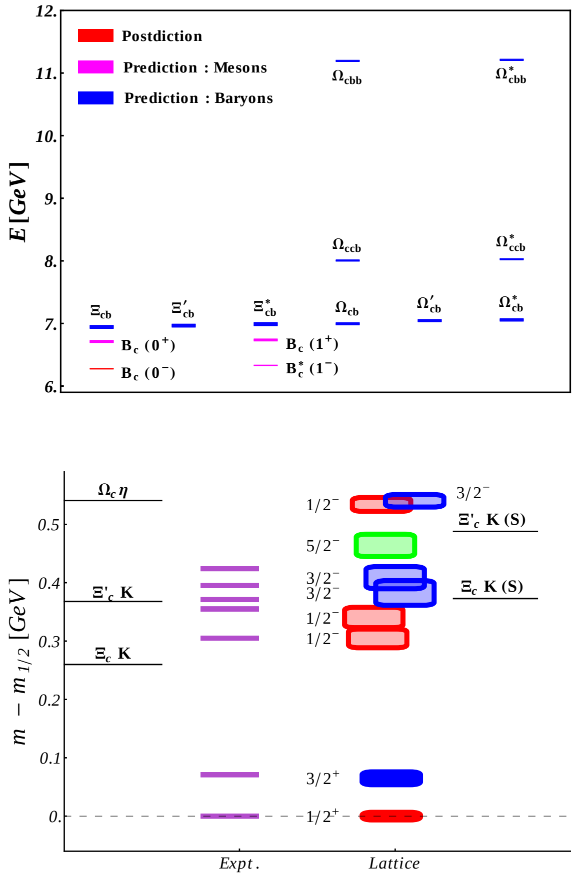



**Fig. 3.** Top: summary of lattice estimates for low lying charmed-bottom hadrons as determined in Ref. [28]. Bottom: Comparison plot from Ref. [32] between the lattice estimates and the experimental values for the energies of $\Omega_c$ excitations.

*Excited baryons* : As discussed in the introduction, one of the major landmark in the year 2017 is the LHCb discovery of five narrow $\Omega_c$ resonances in $\Xi_c^+ K^-$ invariant mass distribution in the energy range between 3000 − 3120 MeV [4]. Following this discovery, Belle collaboration has confirmed four out of these five excited states [5]. Many more highly excited baryons are coming into light with more discoveries. e.g. the observation of a $\Omega^{*-}(3/2^-)$ candidate with a mass of 2012.4(9) MeV by Belle collaboration [29], which is in very good agreement with lattice prediction for such a baryon [30, 31]. Below we discuss the first and only existing lattice investigation of highly excited $\Omega_c$ resonances (Ref. [32]) that predicts the five excited $\Omega_c$ baryons as observed by LHCb.

Following a detailed baryon interpolator construction procedure as invented in Ref. [33,34] a large basis of baryon interpolators, that is expected to extensively scan the radial as well as orbital excitations, are built. By solving the GEVP for correlation matrices constructed out of these interpolators on a lattice ensemble with $m_\pi \sim 391$ MeV (for details see [21]), one extract the $\Omega_c$ baryon spectrum on the lattice. The bottom of Fig. 3 shows a comparison of the lattice energy estimates for the lowest nine $\Omega_c$ excitations with the seven experimentally observed $\Omega_c$ resonances. The relevant strong decay thresholds in the infinite volume are shown as black lines at the top, whereas the black lines at the bottom indicate the relevant non-interacting levels on the lattice. The lowest two levels represent the well known $1/2^+$ and $3/2^+$ excitations. Lattice estimates for these excitations agree well with the experiment. In the energy region, where the five narrow resonances were observed, lattice predicts exactly five levels. Of these five excitations four are in good agreement with the experiment, while the fifth is possibly a $5/2^-$ baryon related to the remaining higher lying experimental candidate. Identifying the quantum information of these lattice levels from the $Z_i^n$s, these five states are argued to be the p-wave excitations [32].

Considering the exploratory nature of this first study, investigating $\Omega_c$ baryon spectrum on multiple lattice ensembles with close to physical $m_\pi$ and larger volumes would be an immediate extension. It would also be an interesting direction to extract the infinite volume scattering matrices considering the allowed baryon-meson scattering channels in the analysis of desired quantum channels in appropriate lattice ensembles. However, the presence of a valence heavy quark, the absence of any valence light quarks and the resonance widths being quite narrow ($< 10$ MeV) [6] indicates our estimates to be robust with such extensive investigations.

## 4 Summary

Over the past decade, lattice QCD has availed multiple precision determinations of the ground state baryon masses using full QCD lattice ensembles with good control over the systematic uncertainties. A summary of lattice determinations of various baryons along with their masses from experiment, where available, are given in Figs. 1, 2 and 3. The only existing exploratory lattice determination of the highly excited $\Omega_c$ states in relation to the recent LHCb discovery and its possible extensions are also discussed.

## 5 Acknowledgements

I would like to thank my collaborators N. Mathur, R. G. Edwards, M. J. Peardon and S. Mondal. I also express my thanks to the organizers and the participants of the Bled workshop for various interesting and insightful discussions. I acknowledge the support from EU under grant no. MSCA-IF-EF-ST-744659 (XQCD-Baryons) and Deutsche Forschungsgemeinschaft Grant No. SFB/TRR 55.

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*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Universal Constituent-Quark Model for All Baryons

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**Abstract.** We discuss the performance of a relativistic constituent-quark model that has been constructed in order to provide a universal framework for the description of all known baryons. After recalling some decisive properties of light-flavor baryons we concentrate on the spectroscopy of baryons containing charm and beauty.

It has become quite evident that modern constituent-quark models serve as an effective tool for accessing quantum chromodynamics (QCD) at low energies. Essential prerequisites are the observations of Poincaré invariance and the spontaneous breaking of chiral symmetry (SB$\chi$S). A model constructed in this spirit has been presented now already 20 years ago [1,2]. It assumes constituent quarks with dynamical masses in a linear confinement according to the string tension of QCD and a hyperfine interaction deriving from Goldstone-boson exchange (GBE), where the latter is cast into pseudoscalar meson exchange. The original version covered all baryons with $u$, $d$, and $s$ flavors. Their spectroscopy is well described in agreement with phenomenology, yielding in particular the inverse level orderings of the first positive- and negative-parity excitations of the nucleon (N), the $J^P = \frac{1}{2}^+$ Roper resonance N(1440) and the $J^P = \frac{1}{2}^-$ N(1535), respectively.

By solving the eigenvalue problem of the pertinent relativistically invariant mass operator one has beyond the eigenvalues (baryon masses) also access to the wave functions of the baryons. Thereby their structures can be tested as revealed under electromagnetic, weak, strong, and gravitational interactions. All of these types of investigations have been carried out with respect to the N over the past. The elastic N electromagnetic form factors calculated in point form - strictly observing Poincaré invariance - have immediately been found in good agreement with experimental data for momentum transfers up to $Q^2 \sim 4$ GeV$^2$ [3,4]. Even the detailed phenomenological insights into the flavor contents of the N electromagnetic form factors are explained correctly, advocating only $\{QQQ\}$ degrees of freedom [5,6].

Similarly, the covariant axial and induced pseudoscalar N form factors have been described in accordance with phenomenology [7,8]. By the same constituent-quark model a microscopic explanation of the strong $\pi NN$ and $\pi N\Delta$ vertex form factors has been provided on the quark level [9]. It largely justifies the phenomenological parameterizations traditionally employed in dynamical models on the hadronic level. Finally the gravitational form factor $A(Q^2)$ of the N has been studied yielding results in accordance with other QCD models, see, e.g., Ref. [10].

This kind of structure studies have also been extended to baryon states other than the N, namely to all of the octet and decuplet baryon ground and some of the resonant states. Of course, in these cases comparisons to experimental data are possible only in a few cases, e.g., for electric radii and magnetic moments [4]. In general, very reasonable results have been found, largely also in good agreement with modern lattice-QCD calculations [11–14].

In view of these results it has come up as an interesting question, if the dynamics of GBE can also be extended to all quark flavors, i.e. all baryon states observed so far. This problem has been answered satisfactorily by the universal relativistic constituent-quark model (URCQM) [10, 15, 16]. It was constructed in the same spirit as its antecessor, the GBE relativistic constituent-quark model (RCQM) of Refs. [1,2], i.e. with the same linear confinement, but now with a pseudoscalar boson exchange of a 24-plet and a singlet, thereby including $u$, $d$, $s$, $c$, and $b$ quark flavors. For the 24-plet GBE only a single mass and a single cut off had to be foreseen with the Goldstone-boson mass equal to the $\pi$ mass and a $\pi$-Q coupling constant derived from the phenomenological $\pi$-N coupling constant using the Goldberger-Treiman relation. The two open parameters inherent in the singlet GBE (the $\eta_0$-Q coupling and the corresponding cut off) were adjusted by fitting the baryon spectra. For more details on the parameter values see Refs. [10] or [16].

Since the 2018 Bled Mini-Workshop concentrated on double-charm baryons, we discuss in this contribution only the description of heavy-flavor baryons by the URCQM [1]. We start with the single-charm spectra in Fig. 1. It is immediately evident that all of the ground states are well reproduced. The same is true with the experimentally established excitations. For the levels in the $\Lambda_c$ and $\Sigma_c$ spectra, where $J^P$ is not definitely known, the URCQM offers nearby levels.

The spectra of the double-charm baryons are shown in Fig. 2. Until recently only for the $\Xi_{cc}$ ground state there were experimental data available from a single experiment, namely the one by the SELEX collaboration [18]. The URCQM produces the $\Xi_{cc}$ ground state more than 100 MeV higher than these data, precisely at 3642 MeV. Similarly, other theoretical models such as the RCQM by the Bonn group [19] (shown by the cyan lines in Fig. 2) or the lattice-QCD calculation by Liu et al. [20] (given by the magenta boxes in Fig. 2) obtain a $\Xi_{cc}$ ground-state level by about the same magnitude higher than the SELEX value.

However, there has been a recent measurement of the $\Xi_{cc}$ ground state by LHCb [21] yielding its mass as $m = 3621.40 \pm 0.72(\text{stat.}) \pm 0.27(\text{syst.}) \pm 0.14(\Lambda_c)$ MeV. The predictions by the URCQM as well as by the other theoretical calculations are now quite compatible with this value. It will be interesting to obtain phenomenological data also for $\Xi_{cc}$ resonances. For the sake of future comparisons we give in Tab. 1 the predictions of the UCRQM for the first seven $\Xi_{cc}$ excitations. We remark that certain $J = \frac{1}{2}$ and $J = \frac{3}{2}$ resonances are degenerate, the reason being that the UCRQM does not (yet) contain tensor forces in the GBE hyperfine interaction. However, as has been learned in case of the GBE RCQM, the inclusion of

[1] The light- and strange-baryon spectra of the URCQM are very similar in most cases identical to the ones of the GBE RCQM of Refs. [1, 2]; cf. also Ref. [15].

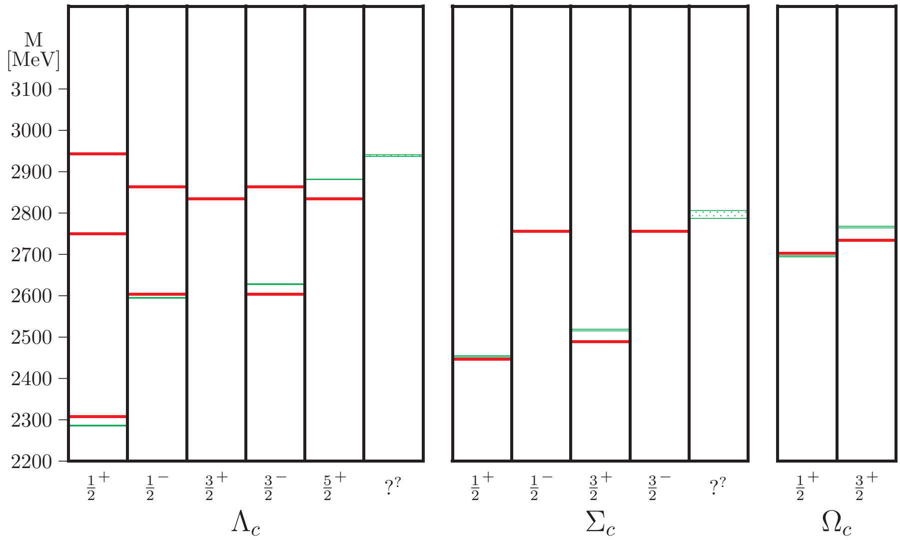


**Fig. 1.** Single-charm baryon spectra of definite $J^P$ as produced by the URCQM (solid/red levels) in comparison to experimental data with their uncertainties (dotted/green levels resp. boxes) [17].

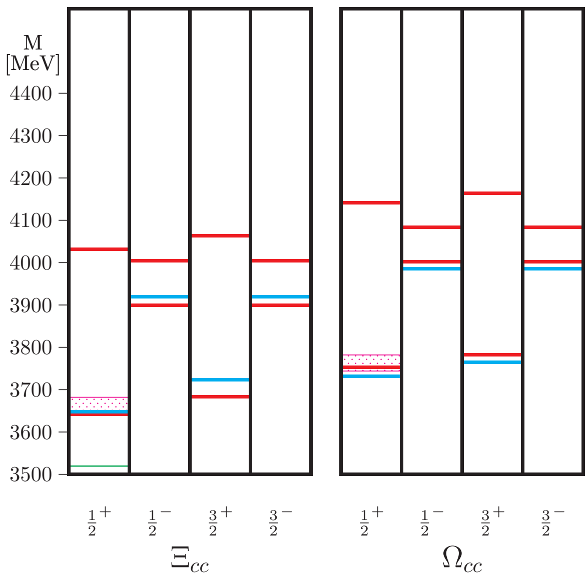


**Fig. 2.** Same as Fig. 1 but for double-charm baryons. The predictions of the URCQM are the solid/red levels. Here they are compared to several other theoretical results and the SELEX experiment for the $\Xi_{cc}$ ground state (lowest/green level) [18]. For further explanations see the text.

all types of hyperfine forces from GBE does not much change the characteristics of the baryon spectra [22,23].

For completeness of the description of heavy baryons by the URCQM we add the single- and double-b baryons in Figs. 3 and 4. Only for single-b baryons we

**Table 1.** $\Xi_{cc}$ ground state and its first seven excitations as predicted by the URCQM.

| State | $J^P$ | URCQM |
|---|---|---|
| $\Xi_{cc}$ | $\frac{1}{2}^+$ | 3642 |
| $\Xi_{cc}$ | $\frac{3}{2}^+$ | 3683 |
| $\Xi_{cc}$ | $\frac{1}{2}^-$ | 3899 |
| $\Xi_{cc}$ | $\frac{3}{2}^-$ | 3899 |
| $\Xi_{cc}$ | $\frac{1}{2}^-$ | 4004 |
| $\Xi_{cc}$ | $\frac{3}{2}^-$ | 4004 |
| $\Xi_{cc}$ | $\frac{1}{2}^+$ | 4032 |
| $\Xi_{cc}$ | $\frac{3}{2}^+$ | 4064 |

may compare to experimental data. In all instances we find good agreement with phenomenology. In case of double-b baryons we are left only with comparisons to other effective models, where we now notice bigger discrepancies.

By the results shown here together with the ones for the light- and strange baryon sectors reported in Refs. [15] and [10] it is certainly evident that a universal relativistic constituent-quark model may be constructed solely on the basis of Goldstone-boson exchange. While beyond the given confinement all of the masses involved in the hyperfine interaction - for constituent quarks u, d, s, c, and b as well as the exchanged Goldstone bosons $\pi$ and $\eta$ - may be taken as predetermined, like the 24-plet coupling constant, there remain only three open fit parameters, the 24-plet as well as the singlet cut offs and the singlet coupling constant. This is certainly remarkable, like the fact that the dynamical ingredients

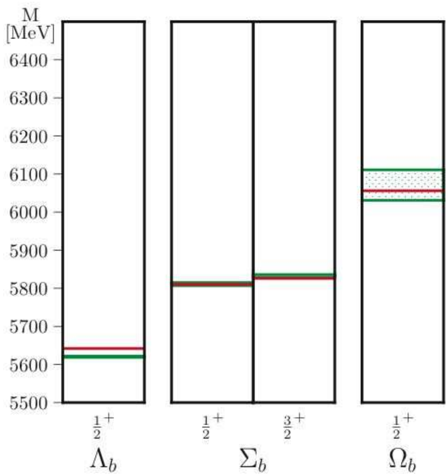


**Fig. 3.** Single-beauty baryon spectra of definite $J^P$ as produced by the URCQM (solid/red levels) in comparison to experimental data with their uncertainties (dotted/green levels resp. boxes) [17].

in the constituent-quark masses (differences between the current and constituent masses) remain practically independent of the flavor [26, 27].

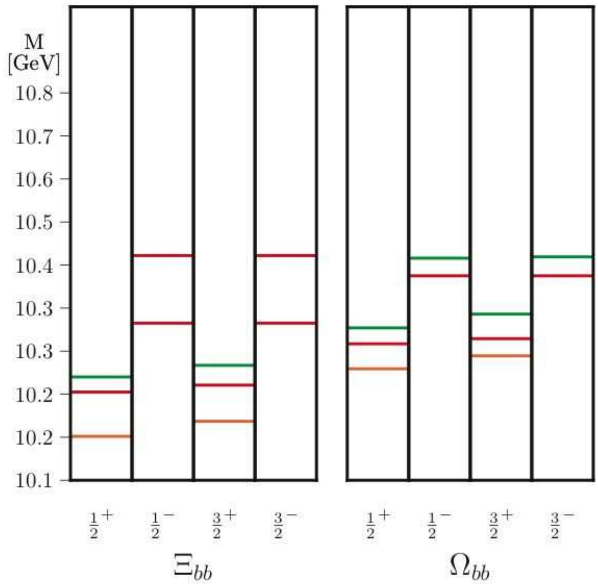


**Fig. 4.** Double-beauty baryon spectra of definite $J^P$ as produced by the URCQM (solid/red levels) in comparison to a nonrelativistic one-gluon-exchange constituent-quark model by Roberts and Pervin [24] (green/higher-lying levels) and the RCQM by Ebert et al. [25] (brown/lower-lying levels).

## Acknowledgement

The author is grateful to Bojan Golli, Mitja Rosina, and Simon Širca for their continuous efforts of organizing every year the Bled Mini-Workshops. These meetings, largely characterized by an informal atmosphere, serve as a valuable institution of exchanging ideas and of mutual learning among an ever growing community of participating colleagues engaged in hadronic physics. In 2018 we have enjoyed the 20th edition thereof, and our best wishes are due for many more such Mini-Workshops to happen in the future.

Results of the GBE relativistic constituent quark model discussed in this contribution have been obtained in collaboration with Ki-Seok Choi and Joseph P. Day.

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BLED WORKSHOPS
IN PHYSICS
VOL. 19, NO. 1
*p. 24*

*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Doubly-heavy baryons, tetraquarks and related topics*

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**Abstract.** We review the physics of doubly-heavy baryons $QQq$ and tetraquarks $QQ\bar{q}\bar{q}$. For the latter, the stability is reached for large enough mass ratio $M/m$, even when spin forces and color mixing are neglected. It is thus customarily claimed that $bb\bar{q}\bar{q}$ in its ground state cannot decay into $b\bar{q} + b\bar{q}$. In some model, $cc\bar{u}\bar{d}$ is shown to be stable if color mixing and spin effects are properly taken into account. It is conjectured that some $bc\bar{q}\bar{q}'$ benefits from favorable adjustments of the gluon tubes in the confinement regime. Some recent studies of pentaquarks and hexaquarks are also summarized.

## 1 Introduction

Double-charm physics, and more generally the physics of doubly-heavy hadrons is by now rather old. Shortly after the prediction of charm by the GIM mechanism [1], Lee, Gaillard and Rosner [2] wrote a seminal paper anticipating many interesting properties of charmed hadrons, including double-charm baryons, with an empirical notation which is now obsolete. As indicated in Sec. 3, the first speculations about $QQ\bar{q}\bar{q}$ arose in 1981 [3], while the first detailed quark model calculation of the doubly heavy baryons $QQq$ came in 1988 [4]. Since then, significant progress has been achieved, with in particular the onset of QCD sum rules and lattice QCD, which is discussed elsewhere in these proceedings. Also, the interaction of light quarks is treated more realistically with the implementation of chiral dynamics. In the abundant literature on $QQq$ and $QQ\bar{q}\bar{q}$, there are also papers with unjustified approximations that do not account for the rich and subtle few-body dynamics inside these hadrons.

## 2 Doubly-heavy baryons

Calculating $QQq$ in a given quark model is rather straightforward, and there are several interesting studies, e.g., [4–6]. In the wave function of the first levels,

* Talk presented by J.-M. Richard

one observes a hierarchy of the average separations, $r(QQ) \ll r(Qq)$, which can be interpreted as a spontaneous or dynamical diquark clustering. But the first excitations occur within the QQ pair and involve a new diquark for each level.

A word about diquarks is in order. There are models where diquarks are introduced as basic constituents. They have a number of successes, and also problems, such as deciding which quarks do cluster in $q_1q_2q_3$, and explaining some nucleon resonances recently seen in photoproduction, which seemingly require two internal excitations. Much more questionable is the diquark picture just as an approximation, or say a lazy way of handling the three-body problem. If the ground-state of $abc$ is searched by solving first the $ab$ problem with the potential $V_{ab}$ alone, and then the two-body problem $[ab]c$ with $V_{ac}(r) + V_{bc}(r)$ with $r$ the distance from $c$ to the center of $[ab]$, then the algebraic energy is underestimated. A simple exercise consists of comparing $V(r)$ and the average of $V(|\mathbf{r} + \mathbf{r}'|)$ when the angles of $\mathbf{r}'$ around $\mathbf{r}$ are varied. Except in the Coulomb case (Gauss theorem), one finds a non-negligible deviation. In other words, the effective QQ interaction within $QQq$ is influenced by the light quark.

Though the Born-Oppenheimer approximation was invented in 1927, it has not yet reached some remote universities. Yet, if any approximation has to be made, this is probably the most interesting. The effective QQ potential in a $QQq$ baryon is the analog of the quark-antiquark potential of charmonium, which itself is also a kind of Born-Oppenheimer potential: the minimal energy of the light degrees of freedom for a given $Q\bar{Q}$ separation.

## 3 Tetraquarks with two heavy quarks

Estimating the tetraquark energy and structure, even in simple quark models, involves a delicate four-body problem. There is a competition between a collective compact configuration and a breaking into two mesons. Unfortunately, this is not always very well handled in the literature. Some authors consistently mistreated the four-body problem in other fields and in quark models. For some other authors, this is more puzzling, as they have set benchmarks of rigor for quarkonium, but became less and less rigorous as the number of constituents was increased. *Corruptio optimorum pessima*[1] use to say our ancestors.

Historically, the first study of $QQ\bar{q}\bar{q}$ was made at CERN [3], with the observation that the system becomes bound, below the $Q\bar{q} + Q\bar{q}$ threshold, if the mass ratio $M/m$ becomes large enough. This was confirmed by Heller et al. [7,8] and Zouzou et al. [9]. The possibility of binding $QQ\bar{q}\bar{q}$ have been rediscovered in some very recent papers, which are sometimes given credit for this idea. This corresponds to the "11th hour effect", *So the last will be first, and the first last* (Matthew 20.16). Another sentence of Matthew's Gospel is also cited in such circumstances, in particular by the sociologist R. Merton [10]: *For to him who has will more be given; and from him who has not, even what he has will be taken away.*

[1] The corruption of the best is the worst

The binding of $QQ\bar{q}\bar{q}$ is a chromoelectric effect at start: the tetraquark benefits from the heavy-heavy attraction that is absent form the threshold. It was also realized that chromomagnetic effects could be decisive for $cc\bar{u}\bar{d}$, with an attraction in the light sector that is absent in the threshold. A decisive progress was accomplished by Janc and Rosina [11], who showed that $cc\bar{u}\bar{d}$ is stable in a specific quark model when chromo-electric and magnetic effects are properly combined. Their result was confirmed and improved by Barnea et al. [12]. See, also, [13].

There are very few rigorous results for the four-body problem, besides the ones shared with any N-body problem, such as the virial theorem and the scaling properties in a power-law potential. The physics of tetraquarks, however, stimulated some contributions: the improved stability when charge-conjugation symmetry C is broken, and the improved stability for asymmetric potentials, as explained below. The first point should have been borrowed from atomic physics, but paradoxically, the quark physics helped to understand the transition from the positronium molecule to the hydrogen one [14,15].[2].

It is well-known that breaking a symmetry lowers the ground-state energy. For instance, going from $H_0 = p^2+x^2$ to $H_0+\lambda\, x$ lowers the first energy from $E_0 = 1$ to $E_0 - \lambda^2/4$, and more generally, breaking parity in $H = H_{\rm even} + H_{\rm odd}$ gives $E < E_{even}$. But in a few-body system, the breaking of symmetry often benefits more to the threshold than to the collective configuration and thus spoils the binding. For instance, in atomic physics, going from $Ps_2$ to $(M^+, m^+, M^-, m^-)$ makes the system unstable for $M/m \gtrsim 2.2$ [17,18]. However, when the symmetry is charge-conjugation, the symmetry breaking benefits entirely to the collective state. Let us, indeed, write the four-body Hamiltonian of the hydrogen molecule as

$$\begin{aligned} H &= \frac{\mathbf{p}_1^2}{2M} + \frac{\mathbf{p}_2^2}{2M} + \frac{\mathbf{p}_3^2}{2m} + \frac{\mathbf{p}_4^2}{2m} + V = H_{\rm even} + H_{\rm odd} \\ &= \left[\sum_i \frac{\mathbf{p}_i^2}{2\mu} + V\right] + \left(\frac{1}{4M} - \frac{1}{4m}\right)\left(\mathbf{p}_1^2 + \mathbf{p}_2^2 - \mathbf{p}_3^2 - \mathbf{p}_4^2\right)\,, \end{aligned} \tag{1}$$

where $2\,\mu^{-1} = M^{-1} + m^{-1}$. The C-parity breaking term, $H_{\rm odd}$, lowers the ground state energy of H with respect to the C-parity even part, $H_{\rm even}$, which is simply a rescaled version of the Hamiltonian of the positronium molecule. Since $H_{\rm even}$ and H have the same threshold, and since the positronium molecule is stable, the hydrogen molecule is even more stable, and stability improves when $M/m$ increases. Clearly, the Coulomb character of V hardly matters in this reasoning, except that if the potential is not Coulombic, $V_{\rm even}$ does not always support a bound state: in this case, stability occurs starting from a minimal value of $M/m$. The key assumption is that the potential does not change when the masses are modified, a property named "flavor independence" in QCD.

As ever, the Born-Oppenheimer approach is very instructive. If one restricts to color $\bar{3}3$, the Born-Oppenheimer QQ potential of $QQ\bar{q}\bar{q}$ is similar to the one of $QQq$, up to an overall constant, which can be identified as the mass difference $Qqq - \bar{Q}q$ from the values at zero separation. See Fig. 1. One thus gets a

[2] To be honest, a similar reasoning was already outlined in the physics of excitons [16]

microscopic derivation of the Eichten-Quigg identity (here without the spin refinements) [19]

$$\mathrm{QQ\bar{q}\bar{q}} \simeq \mathrm{QQq} + \mathrm{Qqq} - \mathrm{\bar{Q}q}\,. \tag{2}$$

Of course, with color mixing, the mass of the tetraquark decreases with respect to the above estimate, and this can be decisive in the charm sector.

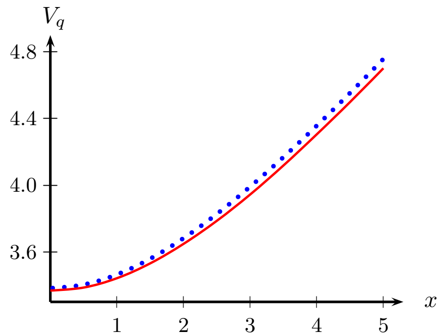


**Fig. 1.** Comparison of the QQ Born-Oppenheimer potentials in QQq (solid line) and $\mathrm{QQ\bar{q}\bar{q}}$ (dotted line), the latter shifted by the mass difference $\mathrm{Qqq} - \mathrm{\bar{Q}q}$

A conservative conclusion, in most studies, is that only $\mathrm{bb\bar{q}\bar{q}}$ is stable. This is indeed the case if spin corrections and color mixing are neglected. With proper inclusion of both color $[\mathrm{QQ}][\mathrm{\bar{q}\bar{q}}] = \bar{3}3$ and $6\bar{6}$ states, and spin effects, one gains some binding in the $\mathrm{cc\bar{u}\bar{d}}$ case. This is shown in Fig. 2.

Another effect could benefit to $\mathrm{bc\bar{q}\bar{q}}$ states. A typical quark model potential reads

$$V = -\frac{3}{16}\sum_{i<j} \tilde{\lambda}_i.\tilde{\lambda}_j \left[-\frac{a}{r_{ij}} + b\, r_{ij} + \frac{\boldsymbol{\sigma}_i.\boldsymbol{\sigma}_j}{m_i\, m_j}\, v_{ss}(r_{ij})\right]\,. \tag{3}$$

The linear part in (3) is interpreted as a string linking the quark to the antiquark. For baryons, it becomes the so-called Y-shape confinement: the three strings join at the Fermat-Torricelli point, to minimize the cumulated length. For a system of two quarks and two antiquarks, a generalization consists of a minimization over the flip-flop and connected double Y arrangements, shown in Fig. 3. The changes with respect to the additive model are minor for baryons, but for tetraquarks, the good surprise is that the flip-flop gives more attraction [20,21], provided the system can evolve freely from one configuration to another one. For identical quarks and/or antiquarks, this is restricted by the Pauli principle. Thus multiquarks with non-identical quarks benefit much better from the string-mediated dynamics. In the future, a comparison of $\mathrm{bb\bar{q}\bar{q}}$, $\mathrm{cc\bar{q}\bar{q}}$ and $\mathrm{bc\bar{q}\bar{q}}$ could probe this effect.

Before leaving the tetraquark sector, let us discuss the all-heavy case $\mathrm{QQ\bar{Q}\bar{Q}}$. It is sometimes claimed to be bound below the $\mathrm{Q\bar{Q}}+\mathrm{Q\bar{Q}}$ threshold, but this is not the case in standard quark models, at least when treated correctly. One may wonder why $\mathrm{Ps}_2$ is demonstrated to be bound [22], and $\mathrm{QQ\bar{Q}\bar{Q}}$ found unstable in a

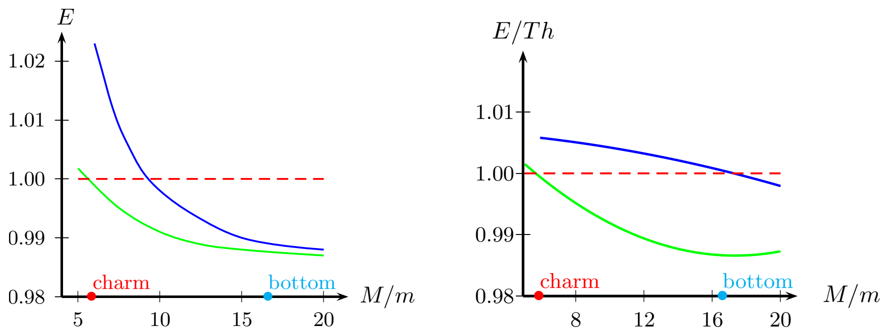


**Fig. 2.** Effect of color-mixing (left) and spin effects (right) on the binding of $QQ\bar{u}\bar{d}$. Left: the tetraquark energy is calculated with only the color $\bar{3}3$ configurations (upper curve) and with the $6\bar{6}$ components (lower curve). Right: the tetraquark energy calculated without (upper curve) or with (lower curve) the chromomagnetic term. The threshold is indicated as a dashed line.

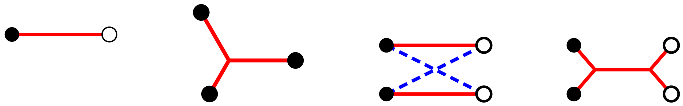

**Fig. 3.** String configurations. From left to right: mesons, baryons, flip-flop and connected double Y for tetraquarks

simple chromoelectric model. Let us consider the generic four-body Hamiltonian

$$H_4 = \sum_i \mathbf{p}_i^2 + \sum_{i<j} g_{ij}\, V(r_{ij})\,, \tag{4}$$

where V is attractive (or dominantly attractive) and $\sum g_{ij} = 2$. For instance, $Ps_2$ corresponds to $V = -1/r$ and $g_{ij} = \{-1, -1, +1, +1, , +1\}$, a tetraquark with color $\bar{3}3$, to $\{1/2, 1/2, 1/4, 1/4, 1/4, 1/4\}$, a tetraquark with color $6\bar{6}$ to $\{-1/4, -1/4, 5/8, 5/8, 5/8, 5/8\}$, and the threshold to $\{1, 1, 0, 0, 0, 0\}$ with a suitable renumbering. The variational principle immediately tells us that the symmetric set of strengths $g_{ij} = 1/3\ \forall i < j$ maximizes the energy, and that increasing the asymmetry of the $g_{ij}$ distribution decreases the energy. The $\chi^2$ of the distribution is larger for $Ps_2$ than for the threshold, and this explains why $Ps_2$ is stable (of course, this is not written exactly in this manner in the textbooks on quantum chemistry!). On the other hand, both color $\bar{3}3$ and $6\bar{6}$ states have a $\chi^2$ smaller than the threshold and thus cannot bind.[3] Numerical calculations show that instability remains when the mixing of color states is accounted for.

[3] To be more precise, if one considers a distribution $\{g_{ij}\} = \{1/3 + 2\lambda, 1/3 + 2\lambda, , 1/3 - \lambda, 1/3 - \lambda, 1/3 - \lambda, 1/3 - \lambda\}$, $E(\lambda_2) < E(\lambda_1)$ is rigorous if $\lambda_2 < \lambda_1 < 0$ or $\lambda_2 > \lambda_1 > 0$, while it is only most plausible if $|\lambda_2| > |\lambda_1|$ with $\lambda_1\,\lambda_2 < 0$, as $E(\lambda)$ is nearly parabolic as a function of $\lambda$.

## 4 Pentaquarks and hexaquarks

Other configurations are regularly revisited, with the hope to predict new stable or metastable multiquarks.

In the pentaquark sector, the $\bar{Q}qqqq$ systems have been revisited. In 1987, it was shown that in the limit where Q is infinitely heavy, and $qqqq = uuds$, $ddus$ or $ssud$ in the $SU(3)_F$ limit, with the assumptions that the strength of the chromomagnetic term is the same as for ordinary baryons, this state is bound by about 150 MeV below the $\bar{Q}q + qqq$ threshold. This pentaquark was searched for in an experiment at Fermilab [23, 24], which turned out inconclusive. The non-strange variant was studied at HERA [25, 26].

More precisely, if $A = \langle -\sum \tilde{\lambda}_i.\tilde{\lambda}_j\, \boldsymbol{\sigma}_i \boldsymbol{\sigma}_j \rangle$ is the expectation value of the chromomagnetic operator for N or $\Lambda$, then $\bar{Q}qqqq$ gets $2\,A$ in the most favorable case. In further studies, it was noticed that as in the case of the famous $H = uuddss$, the multiquark wave function is more dilute than the baryon wave function. This reduces the effectiveness of the chromomagnetic interaction. This is confirmed in our recent study.

Two contributions deal with the hidden-charm states, say $\bar{Q}Qqqq$, which have been much studied after the discovery of the so-called LHCb pentaquarks [27]. First, it is found that within a standard quark model of the type (3), some states are likely below the threshold [28]. This means that new pentaquarks perhaps await discovery, with different quantum numbers.

Another study deals with the states in the continuum. In the early days of the quark model applied to the multiquark sector bound-state techniques were innocently applied to resonances, with the belief that if a state if found, say, 100 MeV above the threshold using a crude one-Gaussian variational wave function, a resonance is predicted at about this energy! The method of real scaling was applied recently to $\bar{c}cuud$ [29], using a standard quark model. It is found that one can separate clearly states that just mimic the continuum from genuine resonances. This is very encouraging, though the candidates for $(3/2)^-$ or $(5/2)^-$ are significantly higher that the LHCb pentaquarks.

In the hexaquark sector, there is a continuous effort from many authors. Our contribution deals with $QQqqqq$, that looks at first very promising, as it combines the chromoelectric attraction of the $QQ$ pair, which acts in the threshold $QQq + qqq$, but not in $Qqq + Qqq$, and the chromomagnetic attraction which is more favorable in the latter than in the former threshold. Moreover, for $qqqq = uuds$, $ddsu$ or $ssud$, the same coherence as in the $\bar{Q}qqqq$ pentaquark could help. However, our study shows that the various effects hardly act together, as each of them requires a specific color-spin configuration.

## 5 Outlook

The physics of multiquark is of primordial importance for hadron spectroscopy. The constituent models, however simple, are a good guidance before considering more ambitious theories. They require some care, but benefit of the know-how accumulated in other branches of few-body physics. Some further developments are required for describing states in the continuum. The method of real scaling looks rather promising, but might be challenged by other schemes. The coupling of channels also reveals interesting features and offers a somewhat complementary point of view [30, 31]. The transition from short-range dynamics in terms of quarks, to a long-range hadron-hadron dynamics is probably the key to describe most of the states.

**Acknowledgments**

JMR would like to thank M. Rosina and his colleagues for the friendly and stimulating atmosphere of the Bled meeting, and also E. Hiyama, M. Oka and A. Hosaka for a fruitlful collaboration about states in the continuum.

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IN PHYSICS
VOL. 19, NO. 1
*p. 32*

*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Production of Pairs of Heavy Quarks by Double Gluon Fusion

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## 1 Double Parton Interactions

The rapid growth of the parton flux at small $x$ gives rise to a dramatic increase of cross sections with large momentum transfer in pp collisions at high energies. In the case of production of mini-jets at the LHC, the inclusive cross section may in fact exceed the value of the total inelastic cross section, for not unrealistically small values of the transverse momenta. One faces therefore a unitarity problem with the large momentum transfer cross sections at high energies, which is solved by introducing Multiple Parton Interactions (MPI) in the process. MPI take into account the possibility of having two or more elementary partonic interactions in a given inelastic hadronic collision and unitarity is restored by MPI because the inclusive cross section is proportional to the multiplicity of interactions. In this way, the inclusive cross section is no more bounded by the value of the total inelastic cross section, when the average multiplicity of interactions is large.

The simplest case of MPI is Double Parton Scattering (DPS). When looking for MPI, one should keep into account that, hard interactions are localised in a space region much smaller as compared to the hadron size and, once the final state is given, the main contribution from MPI is due to the processes which maximise the incoming parton flux.

In DPS the hard component of the interaction is thus disconnected and the non-perturbative components are factorised into functions which depend on two fractional momenta and on the relative transverse distance $b$ between the two interaction points. The non-perturbative input to the DPS cross section, namely the double parton distribution functions, depend therefore explicitly on the relative transverse distance $b$. By neglecting spin and color, the inclusive double parton-scattering cross-section, for two parton processes A and B in a pp collision, is given by [1]:

$$\sigma^D_{(A,B)} = \frac{1}{1+\delta_{A,B}} \sum_{i,j,k,l} \int \Gamma_{i,j}(x_1, x_2; b)\hat{\sigma}^A_{i,k}(x_1, x_1')\hat{\sigma}^B_{j,l}(x_2, x_2')\Gamma_{k,l}(x_1', x_2'; b) \times dx_1\, dx_1'\, dx_2\, dx_2'\, d^2b \quad (1)$$

where the $\Gamma$s represent the double parton distributions and $\hat{\sigma}^{A,B}_{i,k}$ the elementary partonic cross sections. Notice that the dependence of $\sigma^D_{(A,B)}$ on the total transverse energy, of the final state partons with a large $p_t$, is very well characterised and very strong: it is in fact equal to the square of the dependence on the total transverse energy of a single hard scattering cross section. The characteristic dependence of the DPS cross section on the total transverse energy of final state partons with large $p_t$ represents therefore a rather non trivial experimental test of the interaction dynamics.

One may include all unknowns in the process in a quantity with dimensions of a cross section, the *effective cross section*, and the inclusive DPS cross section can thus be expressed by the simplest *pocket formula*, widely used in the experimental analysis of DPS processes:

$$\sigma^D_{(A,B)} = \frac{1}{1+\delta_{A,B}} \frac{\sigma_A \sigma_B}{\sigma_{eff}} \tag{2}$$

where $\sigma_A$ and $\sigma_B$ are the single scattering inclusive pp cross sections for producing the processes A and B respectively. Of course the pocket formula makes sense only if, when comparing with experiment, the effective cross section turns out to be weakly dependent on the kinematics of the process.

The double parton distribution $\Gamma_{i,j}(x_1, x_2; b)$ must therefore depend on the relative transverse distance $b$ between the two partons with fractional momenta $x_1$ and $x_2$. The limiting case of very small $b$ is non-trivial, since in that limit the process cannot be considered a double interaction any more. If $b$ is small enough the two initial state partons can be originated by perturbative splitting, which implies that $\Gamma_{i,j}(x_1, x_2; b)$ is singular in the limit $b \to 0$ [2].

One should thus take into account the contributions to DPS due to parton splitting [3] and that, when splitting is included, some of the contributions to DPS appear also in the single parton scattering cross section [2]. A recent discussion on the matter can be found in [4]: At small $b$, the double parton distributions are conveniently expressed by the sum of a regular and a singular term and the DPS cross section is defined introducing an appropriate regulating function, which cuts the small $b$ region and thus avoids double counting.

On the other hand, for phenomenological studies it is convenient to use the simplest pocket formula, where all unknowns of the process are factorised into a single quantity (the effective cross section). The pocket formula is in fact rather successful in describing the observed DPS cross sections, with values of the effective cross section which show little dependence on the reaction channel and on the kinematical regime (e.g. [5]).

One should stress that the dependence of $\sigma^D_{(A,B)}$ on the total transverse energy of final final state partons is very strong and approximately constant values of $\sigma_{eff}$ represent a non-trivial experimental indication on the interaction process.

## 2 Double $J/\psi$ Production

Double $J/\psi$ production has been studied by several groups, both theoretically and experimentally. A widely used theoretical approach is non-relativistic QCD (NRQCD), an effective theory for heavy quarkonium production (e.g. [7], [8]).

A distinctive feature is that in NRQCD, the heavy quark-antiquark pairs may appear both as color singlet (CS) and color octet (CO) states.

Full LO NRQCD SPS predictions of prompt $J/\psi$, $J/\psi$ hadroproduction (including CS and CO contributions) have been compared [9] with CMS measurements [10]. Although theoretical uncertainties are rather large, SPS at LO underestimates the cross section by more than one order of magnitude, both at small and at large invariant masses and relative rapidities of the $J/\psi$, $J/\psi$ system.

On the other hand, in production of heavy quarks at large $p_t$, the NLO contribution to the cross section can be more than a factor of 10 larger, as compared with the Color Singlet Model contribution at the LO [11]. At large $p_t$, the contribution to the cross section, due to the lowest order diagram, goes in fact as $(1/p_t)^8$, while the contribution to the cross section due to real NLO radiative corrections, has an additional light parton in the final state, which allows the production of a color octet heavy quark pair at a distance of $\mathcal{O}(1/p_t)$. As a consequence, the corresponding contribution to the cross section of heavy quarkonium production goes as $(1/p_t)^6$ at large $p_t$. NLO contributions may thus explain the large difference of the Color Singlet Model LO result with experimental evidence at small invariant masses of the $J/\psi$, $J/\psi$ system.

By working out $J/\psi$ pair production by SPS at 7 TeV c.m. in NRQCD, including NLO contributions, one in fact finds agreement with CMS data at small invariant masses of the $J/\psi$, $J/\psi$ system [12]. The NLO contributions cannot however explain the very strong disagreement (up to three orders of magnitude) of the SPS cross section with CMS data at large invariant masses. To find agreement with the CMS data at large invariant masses one needs to include the contribution of DPS to the process [13].

Double $J/\psi$ production has been studied by LHCb at 13 TeV and the data are compared with theoretical calculations [14]. Also in the LHCb acceptance, SPS at LO gives a negligible contribution to the cross section, except in the case of $k_t$ dependent parton distributions. By evaluating the SPS contribution at the lowest order, with $k_t$ dependent parton distributions, one takes in fact effectively into account most of the NLO contributions.

At small $\Delta Y$ and at small invariant masses, the SPS cross section at NLO is of the right order of magnitude. DPS is on the contrary dominant at large $\Delta Y$ and at large invariant masses. The sum the SPS contribution at NLO, or of SPS at the lowest order with $k_t$ dependent parton distributions, and of the DPS contribution reproduce the data reasonably well.

An extensive study of multiple production of heavy quarks at the LO, using $k_t$ dependent parton distributions and off shell interaction matrix elements, is due to Antoni Szczurek and collaborators [15] [16], which show that in the case of $D^0D^0$ production, DPS dominates also at small invariant masses. The sum of the SPS contribution, at the LO with $k_t$ dependent parton distributions, and of the DPS contribution are able to reproduce the trend of the data LHCb data [17] reasonably well. One should anyhow keep in mind that the $k_t$ dependent parton distributions are still determined with considerable uncertainty and different choices of the $k_t$ dependent parton distributions can easily change the final DPS cross section by a factor two [15].

## 3 Summary

DPS is increasingly important at large c.m. energies and at relatively low transverse momenta and thus play a particularly important role in the case of multiple production of heavy quark pairs at high energies. In the case of production of heavy quarks, the process can in fact be reliably evaluated without introducing cuts in the transferred momenta.

Double $J/\psi$ production has been studied with particular attention both from the experimental and from the theoretical point of view.

The DPS cross sections are characterised by the corresponding values of the "effective cross section" $\sigma_{\mathtt{eff}}$, which although with large experimental uncertainties, turns out to be close to the universal value of 15 mb in almost all measurements. Double $J/\psi$ production by DPS is an exception. In most cases of double $J/\psi$ production by DPS, the observed value of $\sigma_{\mathtt{eff}}$ is in fact smaller (which implies that the production cross section is larger) with respect to the typical values of $\sigma_{\mathtt{eff}}$ observed in all other channels.

At small invariant masses, double $J/\psi$ production is dominated by SPS, which needs to be evaluated at NLO, the LO collinear-factorisation SPS contribution to the cross section being one order of magnitude smaller as compared with presently available data.

At large invariant masses, double $J/\psi$ production is on the contrary dominated (by orders of magnitude) by DPS.

Both for the SPS and for the DPS contributions, a reasonably good agreement with available data on heavy quarks production is obtained by evaluating the cross section at LO with $k_t$ dependent parton distributions and off-shell interaction matrix elements.

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p. 37

*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# The Roper resonance as a meson-baryon molecular state

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**Abstract.** The recently proposed mechanism for the formation of the Roper resonance, in which a dynamically generated state as well as a genuine three-quark resonant state play an equally important role, is confronted with the model proposed almost twenty years ago in which the Roper is pictured as a molecular state of the nucleon and the σ meson.

Our recent investigation on the nature of the Roper resonance [1] has been motivated by the results of lattice QCD simulation in the P11 partial wave by the Graz-Ljubljana and the Adelaide groups [2, 3] that have included beside three-quark interpolating fields also operators for $\pi$N in relative p-wave and $\sigma$N in $s$-wave, and have found no evidence for a dominant three-quark configuration below 1.65 GeV. In our research we use a coupled channel approach which has been previously successfully applied to describe meson scattering and photo- and electro-production in several partial waves in the intermediate energy region [4–9]. In the present analysis of the Roper resonance we include the $\pi$N, $\pi\Delta$, and $\sigma$N channels and solve the Lippmann-Schwinger equation for the meson amplitudes to all orders in the approximation of a separable kernel. We have concluded that while the mass of the resonance is determined by the dynamically generated state, an admixture of the $(1s)^2(2s)^1$ component at an energy around 2 GeV turns out to be crucial to reproduce the experimental width and the modulus of the resonance pole. The mass of the dynamically generated state appears typically 100 MeV below the (nominal) nucleon-sigma threshold. This result agrees well with the prediction of a completely different approach that we studied in the 2001 paper [10] in which we discussed the possibility that the Roper was a molecular state of the nucleon and the $\sigma$ meson. In the following we review the main features of this molecular state and its relation to the dynamically generated state emerging in the coupled channel approach.

In our early approaches to describe the nucleon and the $\Delta(1232)$ we used a chiral version of the linear $\sigma$-model with quarks and determined the quark and meson fields self-consistenly. This model does not work for higher nucleon excitations since the energy of the excited quark turns out to be higher than the free quark mass. In order to ensure confining we used in [10] a chiral version of the Cromodielectric model which included, beside the $\sigma$ and the pion fields, the chromodielectric field $\chi$. The coupling of the $\chi$ field to the quark and meson fields is taken in the form:

$$\mathcal{L}_{\rm int} = \frac{g}{\chi}\,\bar{q}(\hat{\sigma} + i\vec{\tau}\cdot\hat{\vec{\pi}}\gamma_5)q\,, \tag{1}$$

such that for $r \to \infty$, $\chi(r) \to 0$, while the quark mass in this limit behaves as

$$m_q = \frac{g\sigma(r)}{\chi(r)} = \frac{gf_\pi}{\chi(r)} \to \infty\,,$$

which means that the quarks are bound. A typical self-consistent solution for the fields is shown in Fig. 1 a).

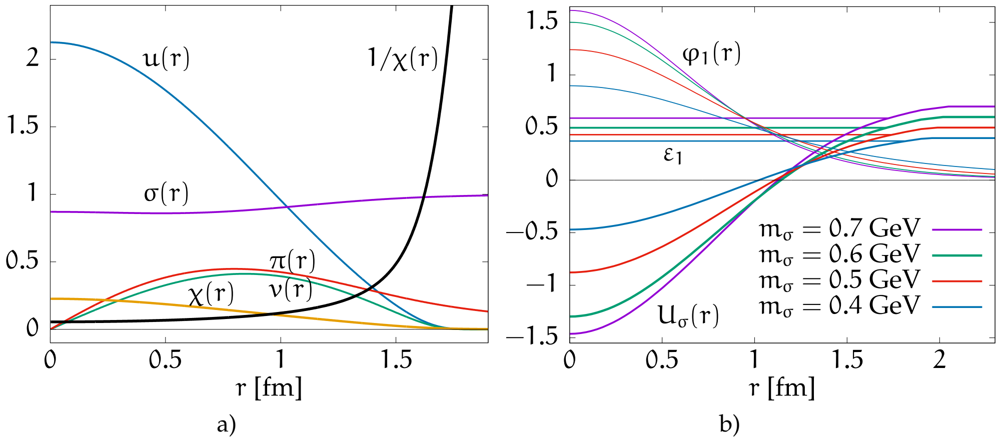


**Fig. 1.** a) Self-consistently determined quark and boson (in units of $f_\pi$) fields in the CDM. b) Effective potential for the $\sigma$ meson and the lowest eigenvalue $\varepsilon_1$ of the corresponding Klein-Gordon equation (in units of GeV) for different choices of the $\sigma$ mass.

We next expanded the field operators of the bosons around their expectation values in the ground state $|N\rangle$; the $\sigma$ operator can be written as:

$$\hat{\sigma}(\mathbf{r}) = \sum_n \frac{1}{\sqrt{2\varepsilon_n}}\,\varphi_n(r)\frac{1}{\sqrt{4\pi}}\left[\tilde{a}_n + \tilde{a}_n^\dagger\right] + \sigma(r)\,, \qquad \tilde{a}_n|N\rangle = 0\,.$$

The stability conditions implies a Klein-Gordon equation for the $\sigma$-meson modes:

$$\left(-\nabla^2 + m_\sigma^2 + U_\sigma(r)\right)\varphi_n(r) = \varepsilon_n^2\varphi_n(r)\,, \qquad U_\sigma(r) = \frac{d^2V(\sigma(r))}{d\sigma(r)^2}\,.$$

Here $V$ stands for the potential originating from (1) and the potential parts of the $\sigma$-model. The potential $U_\sigma$ (see Fig. 1b)) is attractive and supports a bound state which can be interpreted as a molecular state of the nucleon and (one quantum of) the $\sigma$. The corresponding potential for the $\chi$ field turns out to be repulsive, which means that the model does not predict glueball states.

In [10] this excitation of the $\sigma$ field was confronted with the excitation of the quark core in which one quark wass promoted to the 2s orbit. In the self-consistent solution the $2s - 1s$ energy splitting turned out to be smaller than the corresponding vibrational energy $\epsilon_1$, and the conclusion of our work was that the Roper consisted of the dominant quark excitation and a ~ 10 % admixture of the molecular state. However, in that work we used – in accordance with then accepted values – a relatively large $\sigma$ mass between 0.7 GeV and 1.2 GeV. With the present value ~ 0.5 GeV, the lowest eigenmode $\varepsilon_1$ decreases (see Fig. 1b)), while, assuming a somewhat smaller nucleon size, the $2s - 1s$ splitting increases, such that the molecular state may eventually become the dominant component of the Roper resonance.

In our recent paper [1] we study the formation of the resonance in this partial wave in a coupled-channel approach including the $\pi N$, $\pi\Delta$ and $\sigma N$ channels. The Cloudy Bag Model is used to fix the quark-pion vertices while the s-wave $\sigma$-baryon vertex is introduced phenomenologically with the coupling strength $g_\sigma$ as a free parameter and two choices for the mass and the width of the $\sigma$ meson, $m_\sigma = \Gamma_\sigma = 0.6$ GeV and $m_\sigma = \Gamma_\sigma = 0.5$ GeV. Labeling the channels by $\alpha, \beta, \gamma$, the Lippmann-Schwinger equation for the meson amplitude $\chi_{\alpha\gamma}$ for the process $\gamma \to \alpha$ can be cast in the form:

$$\chi_{\alpha\gamma}(k_\alpha, k_\gamma) = \mathcal{K}_{\alpha\gamma}(k_\alpha, k_\gamma) + \sum_\beta \int dk \, \frac{\mathcal{K}_{\alpha\beta}(k_\alpha, k)\chi_{\beta\gamma}(k, k_\gamma)}{\omega(k) + E_\beta(k) - W} \, .$$

Approximating the kernel $\mathcal{K}$ by a separable form, the integral equation reduces to a system of linear equations which can be solved exactly. For sufficiently strong coupling $g_\sigma$ the kernel $\mathcal{K}$ may become singular and a (quasi) bound state arises.

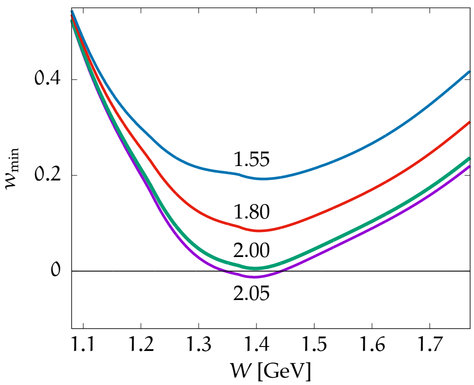


**Fig. 2.** The lowest eigenvalue $w_{\min}$ for four different values of the $\sigma N$ coupling.

| $g_\sigma$ | Re$W_p$ [GeV] | $-2$Im$W_p$ [GeV] |
|---|---|---|
| PDG | 1.370 | 0.175 |
| 1.80 | 1.397 | 0.157 |
| 1.95 | 1.383 | 0.112 |
| 2.00 | 1.358 | 0.111 |
| 2.05 | 1.331 | 0.044 |
| | 1.438 | 0.147 |

**Table 1.** Poles in the complex $W$-plane for four typical values of $g_\sigma$. The PDG values are from [11].

In order to study this process we follow the evolution of the lowest eigenvalue of the matrix pertinent to the system of linear equations, $w_{\min}$, as a function of $W$ for different values of $g_\sigma$ (see Fig. 2). Along with this evolution we observe the evolution of the resonance $S$-matrix pole in the complex $W$-plane using the Laurent-Pietarinen expansion [12–15] (see Table 1). We see that the lowest eigenvalue indeed touches the zero line for $g_\sigma = 2.0$, the pole, however, emerges already for considerably weaker couplings and starts approaching the real axis.

Beyond the critical value, $w_{\min}$ crosses zero twice, producing two poles in the complex energy plane. It is interesting to note that for the values below the critical value, the real part of the pole position almost coincides with $W$ at which $w_{\min}$ reaches its minimum. This value of $W$ is of the order of 100 MeV below the nominal σN threshold. The result agrees well with the molecular picture of the Roper resonance discussed in the first part of this contribution. Let us note that because the σN channel is coupled to other channels, the molecular state has a finite width (i.e. finite $\mathrm{Im}W_\mathrm{p}$) even for $g_s$ greater than the critical value.

In the present approach we have also studied the influence of including a genuine three quark state with one quark excited to the 2$s$ orbit. Using $g_\sigma \approx 1.5$, the results for the position as well as the modulus and the phase come close to the PDG value [11], and are rather insensitive to the mass of the genuine three-quark state. This leads us to the conclusion that the mass of the $S$-matrix pole is determined by the energy of the molecular state while its detailed properties may still considerably depend on the three-quark excited state. The simple model discussed in the first paper provides a simplified picture which enables a deeper insight into the mechanism of the resonance formation, hindered by the complex formalism of the coupled-channel approach.

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*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Double charm baryons and dimesons

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**Abstract.** Several constituent quark models more or less agree in the descripton of baryons and mesons. They may, however, largely disagree in their predictions for dimesons (tetraquarks). The cleanest systems may be ccud = DD* and bbud = (bb)ud. In view of Belle II experiments (at KEK, Japan) in near future, it is of interest to study the DD* dimeson, in order to gain some understanding of the production and detection mechanism and to give some guidance to experimentalists.

## 1 Introduction

There is a strong motivation to verify whether we are capable to extrapolate our experience in QCD from baryons and mesons to many-quark systems. At the level of model-making, it is of interest to look at dimesons (tetraquarks), pentaquarks and dibaryons (hexaquarks). It is important to check whether we may use the same effective quark-quark interaction (apart from the colour factor and the mass-dependent spin-spin term) Vuu = Vcu = Vcc = V$c\bar{c}$ = Vbu = Vbb = V$b\bar{b}$.

The obvious system worth studying is the DD* system which is expected to be either a weakly bound "molecule" or a low-lying resonance. It is relatively long-lived since it decays only electromagnetically (D*$\rightarrow$D$\gamma$) or strongly with extremely small phase space (D*$\rightarrow$D$\pi$). Note that the DD system is a bad candidate since D+D repel each other and no bound state forms.

After the discovery of the $\Xi_{cc}^{++}$ = ccu baryon at LHCb, there is a revived interest for the search of the double charm dimesons. The production mechanism might be similar, but the detection of dimesons is more difficult. For the double charm baryon, they analysed the resonant decay $\Xi_{cc}^{++} \rightarrow \Lambda_c^+ K^- \pi + \pi +$ where the $\Lambda$ baryon was reconstructed in the decay mode $\Lambda_c^+ \rightarrow pK^-\pi^+$. There is no such clear production and detection process available for the DD* intermediate state and it is a great challenge to find a mesurable signature for this dimeson.

## 2 The binding energy of the $\Xi_{cc}^{++}$ baryon

There is a controversy regarding the mass of the double charm baryon. The better documented value of $\Xi_{cc}^{++}$ from LHCb is 3621 MeV while the SELEX value 3519

MeV of $\Xi_{cc}^{+} \to \Lambda_c^{+} K^{-} \pi +$ is met with some scepticism. It would be good to find out whether these are two different states, or SELEX is wrong.

We made a phenomenological estimate using a diquark-quark model and the analogy with mesons (fig.1). Regarding colour quantum number, the diquark in an antisymmetric colour state behaves just as an antiquark. We took a nonrelativistic potential model with a one-gluon-exchange + confining potential with the ”Grenoble parameters AL1” [1] which reproduce rather well most baryons and mesons, in particular also $J/\psi$, the analogon of cc. We get for the mass of the cc diquark 3500 MeV.

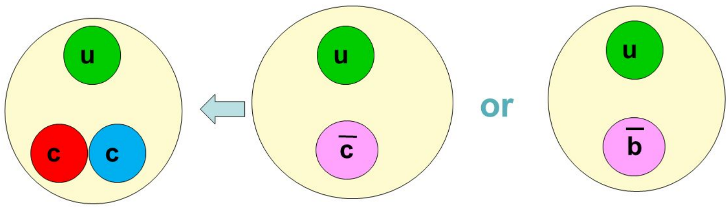


**Fig. 1.** The comparison of the $\Xi_{cc}^{++}$ baryon with the $\bar{D}^0$ and $B^+$ mesons

Using $m(c)$=1870 MeV, $m(\bar{D}^0)$=1865 MeV, $m(b)) = 5259$ MeV and $m(B^+)$=5279 MeV we get

$$m(\Xi_{cc}^{++}) = m(\bar{D}^0) - m(c) + m(cc) = 3495 \text{ MeV}$$

or

$$m(\Xi_{cc}^{++}) = m(B^+) - m(b) + m(cc) = 3520 \text{ MeV}\,.$$

At face value, the latter estimate is very close to the SELEX value. However, the finite size of the diquark and the extra Coulomb repulsion will raise the mass, possibly close to the LHCb value.

Let me quote also other results.

Plessas - the Graz group [2] - obtained with the “Universal constituent quark model for all baryons” (relativistic kinetic energy and a one-Goldston-boson-exchange interaction for the 24-plet + singlet with 5 flavours) the $\Xi_{cc}^{++}$ mass 3642 MeV.

The Lattice QCD result [3] is also around 3600 MeV.

## 3 The binding energy of the D+D* dimeson

In the restricted 4-body space assuming ”cc” in a bound diquark state and the u and d quarks in a general wavefunction, the energy is above the D+D* threshold. In the restricted ”molecular” 4-body space with the two c quarks far apart and a general wavefunction of $\bar{u}$ and $\bar{d}$ the energy is also above the D+D* threshold. Only combining both spaces brings the energy below the threshold.

In the nonrelativistic calculation of Janc and Rosina [4] the one-gluon exchange potential (including the chromomagnetic term) + the linear confining potential was used. The model parameters (Grenoble AL1) [1] fitted all relevant mesons and baryons.

A rich 4-body space was used (an s-state Gaussian expansion at optimized distances, with 3 types of Jacobi coordinates in order to mimic also the p-states. The binding energy (DD*) - (D + D*) = - 2.7 MeV was obtained. This is encouraging, but we have to explore in future, what happens with other interactions and whether the pion cloud between the u and d antiquarks can increase binding, in analogy with the deuteron.

## 4 The formation and decay of the DD* dimeson

There are two possible mechanism for the formation of the dimeson:

1. In the first step the cc-diquark is formed and later automatically dressed by $\bar{u} + \bar{d}$ (or u or d or s in the case of $\Xi_{cc}$ and $\Omega_{cc}$). We have estimated the relative probability of forming ccu or ccd or ccs or the ″atomic″ configuration $cc\bar{u}\bar{d}$ by analogy with the dressing of the b quark into $B^+$, $B^0$, $B^0_s$ and the $\Lambda_b$ baryon determined experimentally in ref. [5]. Initially the relative probability of forming $(cc)\bar{u}\bar{d}$ is about 9% which is about 1/4 probability with respect to $\Xi^{++}_{cc}$ (table 1). Quite a lot! However, this percentage is further reduced by the evolution of the ″atomic″ configuration $(cc)\bar{u}\bar{d}$ into the ″molecular″ configuration of DD*. Mind that the atomic configuration is almost 100 MeV above the D + D* threshold and would decay mostly into two free mesons. The question remains, whether it will decay copiously enough through the DD* bound state or resonance which we are searching for.

**Table 1.** The estimated probability of formation of the tetraquark configuration cc $\bar{u}\bar{d}$

| | | | | | |
|---|---|---|---|---|---|
| $b \to$ | $B^- = b\bar{u}$ | 0.375±0.015 | $cc \to$ | $\Xi^{++}_{cc}$ = ccu | 37% |
| | $B^0 = b\bar{d}$ | 0.375±0.015 | | $\Xi^{+}_{cc}$ = ccd | 37% |
| | $B_s = b\bar{s}$ | 0.160±0.025 | | $\Omega^{+}_{cc}$ = ccs | 16% |
| | $\Lambda_b$ = bud | 0.090±0.028 | | $T^{+}_{cc} = cc\,\bar{u}\bar{d}$ | 9% |

2. In the first step two separate mesons D and D* are formed and then they merge into the DD* dimeson. This process might profit from resonance formation, but due to the dense environement there is a danger that the D + D* system would dissociate before forming the dimeson. The question remains how to distinguish these two mechanisms by analysing the decay products.

The DD* dimeson is stable against a two-body decay into D+D due to its quantum numbers I=0, J=1. It can decay, however, strongly in D+D+$\pi$, or electromagnetically in D+D+$\gamma$, via the decay of D*. The strong decay is very slow (comparable to the electromagnetic decay) due to the extremely small phase space for the pion. Therefore, the DD* dimeson is ”almost stable” and very suitable for detection.

One possibility of detection related to the small phase space of the pionic decay has been proposed by Janc [4, 6]. The ratio between the pionic and gamma decay will strongly depend on the binding or resonance energy of the dimeson. For binding energy more than about 5 MeV there will be only $\gamma$ decay. But there will be a strong background due to the decay of free B* and some kinematical analysis is needed to distinguish it.

## 5 Conclusion

More work is needed to predict theoretically the mechanism of formation of the DD* dimeson and to suggest to experimentalists a reliable signature or tagging.

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IN PHYSICS
VOL. 19, NO. 1
*p. 45*

*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Pion electroproduction in the energy region of the Roper resonance

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The most detailed and model-independent experimental studies of the structure of the Roper resonance utilize coincident electron scattering, in particular in the neutral-pion electroproduction channel. The most recent such study, exploiting polarization degrees of freedom to enhance the sensitivity to the pertinent electroproduction multipoles, has been performed at the three spectrometer facility of the A1 Collaboration at the Mainz Microtron (MAMI) [1]. The $p(\vec{e}, e'\vec{p})\pi^0$ process has been investigated at $W \approx (1440 \pm 40)$ MeV, $Q^2 \approx (0.1 \pm 0.02)(\mathrm{GeV}/c)^2$ and $\theta_p^* \approx (90 \pm 15)^\circ$.

Two helicity-dependent recoil polarization components, $P_x'$ and $P_z'$, have been extracted, as well as the helicity-independent component $P_y$. and compared to the values calculated by MAID [2], DMT [3] and SAID [4]. No model reproduces all features of the data simultaneously. The scalar helicity amplitude $S_{1/2}$ — sensitive to potentially large pion-cloud effects at such small $Q^2$ — has also been determined. Fixing the transverse helicity amplitude $A_{1/2}$ to its MAID value and taking $S_{1/2}^{\mathrm{MAID}}$ as the nominal best model value, we have been able to express $S_{1/2}$ from our fit as the fraction of $S_{1/2}^{\mathrm{MAID}}$, resulting in

$$S_{1/2} = \left(0.80^{+0.15}_{-0.20}\right) S_{1/2}^{\mathrm{MAID}} = \left(14.1^{+2.6}_{-3.5}\right) \cdot 10^{-3}\mathrm{GeV}^{-1/2} .$$

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*Proceedings of the Mini-Workshop*
***Double-charm baryons and dimesons***
Bled, Slovenia, June 17 - 23, 2018

# Discussion sessions

Workshop participants

At the Mini-workshop Bled 2018, there were also two discussion sessions with informal presentations and discussions.

Bogdan Povh (Heidelberg) described new ideas how to explain the well-known deficit in the nucleon spin polarization. It can be better understood if the $x$-dependence of the polarized structure functions is analysed and not only their integral. He found that the measured structure function agrees with the prediction of the static quark model for Bjorken variable $x > 0.2$ and drops rapidly for $x < 0.2$. The interpretation is that for $x > 0.2$ electrons get scattered quasi-elastically on the undamaged constituent quarks and for $x < 0.2$ on the fragmented constituent quarks which do not preserve the polarization. In Fig. 1 the $x$-dependence of the static constituent quark model is represented by the unpolarized structure function $F^p(x)$ (full line).

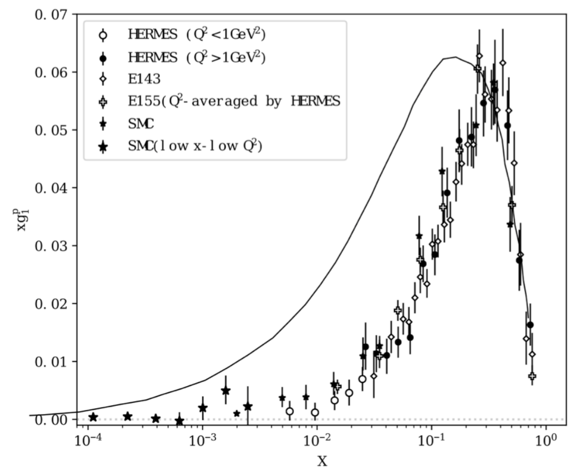


**Fig. 1.** Comparison of the prediction of the statical model (full line) and the data for the polarized structure function.

Full derivation was presented later at the workshop on Diffraction and Low-$x$ in Reggio Calabria (Aug. 26 - Sept. 1, 2018) [1].

The second discussion session was devoted to the news from Lattice QCD. Saša Prelovšek (Ljubljana) reported her work on QCD at high temepartures and its symmetries from a lattice study [2]. The observed degeneracies in the spectrum are similar to Glozman's depleted spectrum at T=0. The additional chiral-spin symmetry $SU(2)_{CS}$ appears, as well as a SU(4) symmetry in the limit of high T, even if these symmetries are not present in the Lagrangian. The $SU(2)_{CS}$ symmetry manifests itself in the degeneracy of the vector and tensor-vector spatial correlators as the temperature is increased to about $T = 380$ MeV (Fig. 2 b)). At this temperature the ratio of these correlators, that are not related by the chiral transformation, approaches one.

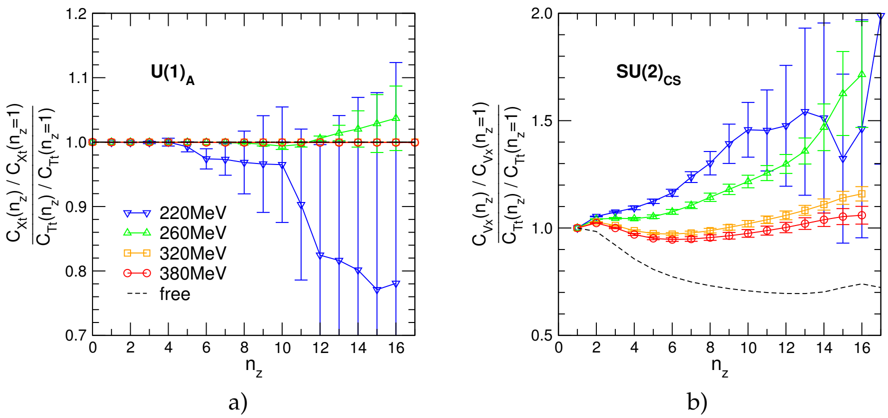


**Fig. 2.** Ratios of normalized correlators, that are related by $U(1)_A$ and $SU(2)_{CS}$ symmetry.

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# Povzetki v slovenščini

## Mezonski spektri: eksperimentalni podatki in njihovo tolmačenje

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Pomembno je pravilno tolmačenje eksperimentalnih podatkov. Predvsem poudarjamo, da bi morali spektre sistemov kvark-antikvark proučevati iz konfiguracij z dobrimi kvantnimi števili, pri tem pa je najprimernejši sistem čarmonij. Predlagamo možna bodoča odkritja na podlagi obstoječih slabih podatkov o čarmoniju in botomoniju. Pregledamo tudi naša zapažanja glede bozonov E(38 MeV) in Z(57.5 GeV).

## Spektroskopija težkih barionov v kromodinamiki na mreži

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V tem poročilu obravnavam najnovejše in najnatančnejše ocene mas barionov v osnovnem stanju z uporabo kromodinamike na mreži. Glede na dobre izglede za področje težkih barionov poudarjam zlasti ocene zanje. Obravnavam tudi prvo in edino določitev visokih vzbujenih stanj bariona $\Omega_c$ na mreži v povezavi z nedavnim odkritjem na LHCb.

## Univerzalni model iz konstituentnih kvarkov za vse barione

W. Plessas

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Obravnavamo uspešno delovanje relativističnega modela iz konstituentnih kvarkov, ki smo ga izdelali kot splošno orodje za opis vseh znanih barionov. Najprej ponovimo nekaj odločilnih značilnosti lahkih barionov, potem pa se osredotočimo na spektroskopijo barionov, ki vsebujejo kvarke c in b.

## Dvojno težki barioni in tetrakvarki ter sorodni sistemi

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Podam pregled fizike dvojno težkih barionov $QQq$ in tetrakvarkov $QQ\bar{q}\bar{q}$. Te zadnji postanejo stabilni pri dovolj velikem razmerju mas $M/m$ , tudi če zanemarimo spinske sile in mešanje barv. Običajno zatrjujejo, da $bb\bar{q}\bar{q}$ v osnovnem stanju ne more razpasti v $b\bar{q}$ + $b\bar{q}$. V nekaterih modelih pa so pokazali, da je tudi $cc\bar{q}\bar{q}$ stabilen, če pravilno upoštevamo mešanje barv in spinske efekte. Domnevamo, da nekatera stanja $bc\bar{q}\bar{q}'$ lahko izkoristijo ugodne prilagoditve cevi gluonov v področju konfinacije. Omenim tudi nekatere nedavne raziskave pentakvarkov in heksakvarkov.

## Proizvodnja parov težkih kvarkov z dvojnim gluonskim zlivanjem

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Hiter porast toka partonov z majhnim deležem gibalne količine $x$ povzroči dramatično povečanje preseka pri trkih visokoenergijskih protonov na protonih. Inkluzivni presek za mini-pljuske lahko preseže celotni neelastični presek. Problem unitarnosti reši vpeljava večpartonske (vsaj dvopartonske) interakcije. Vpeljemo kot parameter *efektivni presek*, ki vključuje vse neznane količine pri procesu, in uspešno razložimo dvojno produkcijo mezonov $J/\psi$. Pri majhnih invariantnih masah prevledujejo enojni trki partonov, pri velikih invariantnih masah pa dvojni.

## Roperjeva resonanca kot molekulsko stanje bariona in mezona

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Mehanizem, ki smo ga pred kratkim predlagali za opis nastanka Roperjeve resonance, pri katerem igrata enako pomembni vlogi dinamično generirano stanje in vzbujeno stanje treh kvarkov, konfrontiramo z modelom, predlaganim pred skoraj dvajsetimi leti, v katerem Roperjevo resonanco obravnavamo kot molekulo, sestavljeno iz nukleona in mezona sigma.

## Dvojno čarobni barioni in dimezoni

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Mnogi modeli iz konstituentnih kvarkov se bolj ali manj ujemajo pri opisu mezonov in barionov. Lahko se pa močno razlikujejo pri napovedih za dimezone (tetrakvarke). Najčistejša sistema utegneta biti $cc\bar{u}\bar{d}$ = DD* in $bb\bar{u}\bar{d}$ = (bb)$\bar{u}\bar{d}$. Glede na bodoče meritve z detektorjem Belle II na Japonskem se izplača proučevati dimezone DD*, da bi bolje razumeli mehanizem proizvodnje in prepoznavanja in bi s tem dali nekaj opore eksperimentalcem.

## Pionska elektroprodukcija v energijskem področju Roperjeve resonance

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Roperjeva resonanca in njena elektromagnetna struktura sodita med pomembne nerešene uganke sodobne hadronske fizike. Lastnosti tega najnižjega vzbujenega stanja nukleona z istimi kvantnimi števili so težko dostopne, saj je resonanca skrita pod velikim ozadjem sosednjih resonanc. V prispevku smo poročali o meritvi polarizacijskih komponent odrinjenega protona iz procesa $p(\vec{e}, e'\vec{p})\pi^0$, in sicer od vijačnosti odvisnih $P'_x$, $P'_z$ ter od vijačnosti neodvisne $P_y$. Rezultate smo primerjali z modelskimi izračuni MAID, DMT in SAID ter ugotovili neujemanje zlasti pri slednjih dveh. Ob določenih modelskih privzetkih smo določili tudi skalarno vijačnostno amplitudo $S_{1/2}$.

## Diskusije

Udeleženci delavnice

Imeli smo tudi dve srečanji z neformalnimi predstavitvami in diskusijami. Bogdan Povh (Heidelberg je predstavil novo zamisel, kako razložiti znani primanjkljaj pri polarizaciji nukleonovega spina. Če pogledamo odvisnost polarizirane strukturne funkcije v odvisnosti od deleža gibalne količine $x$ (in ne, kot običajno, le njenega integrala), vidimo, da nastopi primanjkljaj pri majhnih $x$. To pojasnimo, da se elektroni pri $x > 0,2$ sipljejo na celem kvarku, pri $x < 0,2$ pa na "razbitinah" kvarka.

Saša Prelovšek (Ljubljana) pa je poročala o svojem delu na kromodinamiki na mreži pri visokih temperaturah. Pojavi se degeneracija v spektru, ki kaže na dodatni simetriji $SU(2)_{CS}$ in SU(4), ki nista prisotni v Lagrangevi funkciji.

Blejske Delavnice Iz Fizike, Letnik 19, št. 1, ISSN 1580-4992

Bled Workshops in Physics, Vol. 19, No. 1

Zbornik delavnice 'Dvojno čarobni barioni in dimezoni',
Bled, 17. – 23. junij 2018

Proceedings of the Mini-Workshop 'Double-charm baryons and dimesons',
Bled, June 17 – 23, 2018

Uredili in oblikovali Bojan Golli, Mitja Rosina, Simon Širca

Članki so recenzirani. Recenzijo je opravil uredniški odbor.

Izid publikacije je finančno podprla Javna agencija za raziskovalno dejavnost RS iz sredstev državnega proračuna iz naslova razpisa za sofinanciranje domačih znanstvenih periodičnih publikacij.

Tehnični urednik Matjaž Zaveršnik

Založilo: DMFA – založništvo, Jadranska 19, 1000 Ljubljana, Slovenija

Natisnila tiskarna Itagraf v nakladi 65 izvodov

Publikacija DMFA številka 2076

Brezplačni izvod za udeležence delavnice