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Mass spectra and Regge Trajectories of Heavy Mesons and Baryons

Dietmar Ebert*
Institute of Physics, Humboldt University, Germany
E-mail: debert@physik.hu-berlin.de

Rudolf N. Faustov
Dorodnicyn Computing Centre, Russian Academy of Sciences, Vavilov Str. 40, 119991 Moscow, Russia
E-mail: faustov@ccas.ru

Vladimir O. Galkin
Dorodnicyn Computing Centre, Russian Academy of Sciences, Vavilov Str. 40, 119991 Moscow, Russia
E-mail: galkin@ccas.ru

Masses of the ground state, orbitally and radially excited heavy–light mesons and heavy baryons are calculated within the framework of the relativistic quark model based on the quasipotential approach. Heavy baryons are considered in the heavy-quark–light-diquark picture. The Regge trajectories of heavy mesons and baryons for orbital and radial excitations are constructed, and their linearity, parallelism, and equidistance are verified.

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*Speaker.

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1. Introduction

Recently significant experimental progress has been achieved in studying the spectroscopy of hadrons with one heavy ($Q = c, b$) quark [1]. In the meson sector several new excited states of heavy-light mesons were discovered, some of which have rather unexpected properties [1]. The most investigated and intriguing issue is the charmed-strange meson sector, where masses of nine mesons have been measured. Even eight years after the discovery of $D_{s0}^*(2317)$ and $D_{s1}(2460)$ mesons their nature remains controversial in the literature. The abnormally light masses of these mesons put them below DK and D^*K thresholds thus making these states narrow since the only allowed decays violate isospin. The peculiar feature of these mesons is that they have masses either almost equal or even lower than the masses of their charmed counterparts $D_0^*(2400)$ and $D_1(2427)$. Most of the theoretical approaches including lattice QCD, QCD sum rules and different quark model calculations yield masses of the 0^+ and 1^+ P-wave $c\bar{s}$ states significantly heavier (by 100-200 MeV) than the measured ones. Different theoretical solutions of this problem were proposed including consideration of these mesons as chiral partners of 0^- and 1^- states , $c\bar{s}$ states which are strongly influenced by the nearby DK thresholds, DK or $D_s\pi$ molecules, a mixture of $c\bar{s}$ and tetraquark states. However the clear understanding of their nature is still missing.

In the baryon sector the number of the observed charmed and bottom baryons almost doubled in last few years and now it is nearly the same as the number of known charmed and bottom mesons. Observations of new charmed baryons were mainly done at the B-factories, while new bottom baryons were discovered at Tevatron. It is expected that new data on excited bottom baryons will come soon from the LHC, where they are supposed to be copiously produced. Due to the poor statistics, the quantum numbers of most of the excited states of heavy baryons are not known experimentally and are usually prescribed following the quark model predictions [1].

In this talk we will briefly review our recent investigations of heavy meson and baryon spectroscopy in the framework of the QCD-motivated relativistic quark model based on the quasipotential approach. The dynamics of all quarks in hadrons is treated completely relativistically without application of either the nonrelativistic v/c or heavy-quark $1/m_Q$ expansions. To simplify the very complicated relativistic three-body problem heavy baryons are considered in the heavy-quark–light-diquark approximation. Such scheme significantly reduces the number of the excited baryon states compared to the genuine three-quark picture. Here we present the results of the calculation of the masses of the excited heavy mesons and baryons up to rather high orbital and radial excitations. This allows us to construct the heavy meson and baryon Regge trajectories both in the (J, M^2) and (n_r, M^2) planes, where J is the hadron spin, M is the hadron mass and n_r is the radial quantum number. Then we can test their linearity, parallelism and equidistance and determine their parameters: slopes and intercepts. Their properties are of great importance, since they provide a better understanding of the hadron structure. Moreover, their knowledge is also important for non-spectroscopic problems such as, e.g., hadron production and high energy scattering.

2. Relativistic quark model

In the relativistic quark model based on the quasipotential approach a meson or a baryon in the quark-diquark picture is described by the wave function of the corresponding bound state, which

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satisfies the quasipotential equation of the Schrödinger type [2]

$$\left(\frac{b^2(M)}{2\mu_R}-\frac{\mathbf{p}^2}{2\mu_R}\right)\Psi_M(\mathbf{p})=\int\frac{d^3q}{(2\pi)^3}V(\mathbf{p},\mathbf{q};M)\Psi_M(\mathbf{q}), \tag{2.1}$$

where the relativistic reduced mass is

$$\mu_R=\frac{E_1E_2}{E_1+E_2}=\frac{M^4-(m_1^2-m_2^2)^2}{4M^3}, \tag{2.2}$$

and E_1, E_2 are given by

$$E_1=\frac{M^2-m_2^2+m_1^2}{2M},\quad E_2=\frac{M^2-m_1^2+m_2^2}{2M}. \tag{2.3}$$

Here $M=E_1+E_2$ is the bound state mass (meson, diquark or baryon), $m_{1,2}$ are the quark (diquark) masses, and $\mathbf{p}$ is their relative momentum. In the center-of-mass system the relative momentum squared on mass shell reads

$$b^2(M)=\frac{[M^2-(m_1+m_2)^2][M^2-(m_1-m_2)^2]}{4M^2}. \tag{2.4}$$

The kernel $V(\mathbf{p},\mathbf{q};M)$ in Eq. (2.1) is the quasipotential operator of the quark (diquark) interaction. It is constructed with the help of the off-mass-shell scattering amplitude, projected onto the positive energy states. Constructing the QCD-motivated quasipotential of the interquark interaction, we have assumed that the effective interaction is the sum of the usual one-gluon exchange term with the mixture of long-range vector and scalar linear confining potentials, where the vector confining potential contains the Pauli interaction. The quasipotential is then defined by
(a) quark-antiquark ($q\bar{Q}$) interaction (meson)

$$V_{q\bar{Q}}(\mathbf{p},\mathbf{q};M)=\bar{u}_1(p)\bar{u}_2(-p)\mathcal{V}(\mathbf{p},\mathbf{q};M)u_1(q)u_2(-q), \tag{2.5}$$

with

$$\mathcal{V}(\mathbf{p},\mathbf{q};M)=\frac{4}{3}\alpha_sD_{\mu\nu}(\mathbf{k})\gamma_1^\mu\gamma_2^\nu+V^V_{\rm conf}(\mathbf{k})\Gamma_1^\mu\Gamma_{2;\mu}+V^S_{\rm conf}(\mathbf{k}),$$

(b) quark-quark (qq) interaction in the colour antitriplet state (diquark)

$$V_{qq}(\mathbf{p},\mathbf{q};M)=\bar{u}_1(p)\bar{u}_2(-p)\mathcal{V}(\mathbf{p},\mathbf{q};M)u_1(q)u_2(-q), \tag{2.6}$$

with

$$\mathcal{V}(\mathbf{p},\mathbf{q};M)=\frac{1}{2}\left[\frac{4}{3}\alpha_sD_{\mu\nu}(\mathbf{k})\gamma_1^\mu\gamma_2^\nu+V^V_{\rm conf}(\mathbf{k})\Gamma_1^\mu(\mathbf{k})\Gamma_{2;\mu}(-\mathbf{k})+V^S_{\rm conf}(\mathbf{k})\right],$$

(c) quark-diquark (Qd) interaction in the colour singlet state (baryon)

$$\begin{aligned}V_{Qd}(\mathbf{p},\mathbf{q};M) &= \frac{\langle d(P)|J_\mu|d(Q)\rangle}{2\sqrt{E_d(p)E_d(q)}}\bar{u}_Q(p)\frac{4}{3}\alpha_sD_{\mu\nu}(\mathbf{k})\gamma^\nu u_Q(q)\\ &\quad+\psi_d^*(P)\bar{u}_Q(p)J_{d;\mu}\Gamma_Q^\mu(\mathbf{k})V^V_{\rm conf}(\mathbf{k})u_Q(q)\psi_d(Q)\\ &\quad+\psi_d^*(P)\bar{u}_Q(p)V^S_{\rm conf}(\mathbf{k})u_Q(q)\psi_d(Q),\end{aligned} \tag{2.7}$$

where α_s is the QCD coupling constant, $D_{\mu\nu}$ is the gluon propagator in the Coulomb gauge, and $\mathbf{k}=\mathbf{p}-\mathbf{q}$; γ_μ and $u(p)$ are the Dirac matrices and spinors, $\langle d(P)|J_\mu|d(Q)\rangle$ is the vertex of the

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diquark-gluon interaction which takes into account the internal diquark structure, $P=(E_d(p),-\mathbf{p})$, $Q=(E_d(q),-\mathbf{q})$ and $E_d(p)=\sqrt{\mathbf{p}^2+M_d^2}$.

The effective long-range vector vertex of the quark is given by

$$\Gamma_\mu(\mathbf{k})=\gamma_\mu+\frac{i\kappa}{2m}\sigma_{\mu\nu}k^\nu, \tag{2.8}$$

where κ is the Pauli interaction constant characterizing the anomalous chromomagnetic moment of quarks. Vector and scalar confining potentials in the nonrelativistic limit reduce to

$$\begin{aligned} V^V_{\rm conf}(r) &= (1-\varepsilon)(Ar+B),\\ V^S_{\rm conf}(r) &= \varepsilon(Ar+B), \end{aligned} \tag{2.9}$$

reproducing

$$V_{\rm conf}(r)=V^S_{\rm conf}(r)+V^V_{\rm conf}(r)=Ar+B, \tag{2.10}$$

where ε is the mixing coefficient. The diquark state in the confining part of the quark-diquark quasipotential (2.7) is described by the wave functions

$$\psi_d(p)=\begin{cases} 1 & \text{for the scalar diquark}\\ \varepsilon_d(p) & \text{for the axial vector diquark}\end{cases}, \tag{2.11}$$

where ε_d is the polarization vector of the axial vector diquark. The effective long-range vector vertex of the diquark can be presented in the form

$$J_{d;\mu}=\begin{cases} \dfrac{(P+Q)_\mu}{2\sqrt{E_d(p)E_d(q)}} & \text{for the scalar diquark}\\ -\dfrac{(P+Q)_\mu}{2\sqrt{E_d(p)E_d(q)}}+\dfrac{i\mu_d}{2M_d}\Sigma_\mu^\nu\tilde{k}_\nu & \text{for the axial vector diquark}\end{cases}, \tag{2.12}$$

where $\tilde{k}=(0,\mathbf{k})$. Here Σ^ν_μ is the antisymmetric tensor

$$\left(\Sigma_{\rho\sigma}\right)^\nu_\mu=-i(g_{\mu\rho}\delta^\nu_\sigma-g_{\mu\sigma}\delta^\nu_\rho), \tag{2.13}$$

and the axial vector diquark spin $\mathbf{S}_d$ is given by $(S_{d;k})_{il}=-i\varepsilon_{kil}$. We choose the total chromomagnetic moment of the axial vector diquark $\mu_d=0$.

The constituent quark masses $m_u=m_d=0.33$ GeV, $m_s=0.5$ GeV, $m_c=1.55$ GeV, $m_b=4.88$ GeV, and the parameters of the linear potential $A=0.18$ GeV^2 and $B=-0.3$ GeV were fixed in our previous calculations. The value of the mixing coefficient of vector and scalar confining potentials $\varepsilon=-1$ has been determined from the consideration of charmonium radiative decays and the heavy quark expansion. Finally, the universal Pauli interaction constant $\kappa=-1$ has been fixed from the analysis of the fine splitting of heavy quarkonia ^3P_J- states. Note that the long-range chromomagnetic contribution to the potential in our model is proportional to $(1+\kappa)$ and thus vanishes for the chosen value of $\kappa=-1$.

Since we deal with mesons, diquarks and baryons containing light quarks we adopt for the QCD coupling constant $\alpha_s(\mu^2)$ the simplest model with freezing, namely

$$\alpha_s(\mu^2)=\frac{4\pi}{\beta_0\ln\dfrac{\mu^2+M_B^2}{\Lambda^2}}, \qquad \beta_0=11-\frac{2}{3}n_f, \tag{2.14}$$

where the scale is taken as $\mu=2m_1m_2/(m_1+m_2)$, the background mass is $M_B=2.24\sqrt{A}=0.95$ GeV, and $\Lambda=413$ MeV was fixed from fitting the ρ mass [3].

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3. Heavy-light mesons

The quasipotential (2.5) can in principal be used for arbitrary quark masses. The substitution of the Dirac spinors into (2.5) results in an extremely nonlocal potential in the configuration space. Clearly, it is very hard to deal with such potentials without any additional approximations. In order to simplify the relativistic $q\bar{Q}$ potential, we make the following replacements in the Dirac spinors:

$$\varepsilon_{1,2}(p) = \sqrt{m_{1,2}^2 + \mathbf{p}^2} \to E_{1,2} \tag{3.1}$$

(see the discussion of this point in [3]). This substitution makes the Fourier transform of the potential (2.5) local. The resulting $q\bar{Q}$ potential then reads

$$V_{q\bar{Q}}(r) = V_{\rm SI}(r) + V_{\rm SD}(r), \tag{3.2}$$

where the explicit expression for the spin-independent $V_{\rm SI}(r)$ can be found in Ref. [3]. The structure of the spin-dependent potential is given by

$$V_{\rm SD}(r) = a_1\,\mathbf{L}\mathbf{S}_1 + a_2\,\mathbf{L}\mathbf{S}_2 + b\left[-\mathbf{S}_1\mathbf{S}_2 + \frac{3}{r^2}(\mathbf{S}_1\mathbf{r})(\mathbf{S}_2\mathbf{r})\right] + c\,\mathbf{S}_1\mathbf{S}_2 + d\,(\mathbf{L}\mathbf{S}_1)(\mathbf{L}\mathbf{S}_2), \tag{3.3}$$

where $\mathbf{L}$ is the orbital angular momentum, $\mathbf{S}_i$ is the quark spin. The coefficients a_1, a_2, b, c and d are expressed through the corresponding derivatives of the Coulomb and confining potentials. Their explicit expressions are given in Ref. [3].

The calculated masses of heavy-light D and D_s mesons are given in Table 1 ($n = n_r + 1$, L is the orbital momentum, J and S are the total angular momentum and spin). They are confronted with available experimental data from PDG [1]. The results for B and B_s mesons are given in Ref. [4].

In our analysis we calculated masses of both orbitally and radially excited heavy-light mesons up to rather high excitation numbers ($L = 4$ and $n_r = 4$). This makes it possible to construct the heavy-light meson Regge trajectories in the (J, M^2) and (n_r, M^2) planes. We use the following definitions.

(a) The (J, M^2) Regge trajectory:

$$J = \alpha M^2 + \alpha_0; \tag{3.4}$$

(b) The (n_r, M^2) Regge trajectory:

$$n_r = \beta M^2 + \beta_0, \tag{3.5}$$

where α, β are the slopes and α_0, β_0 are intercepts. The relations (3.4) and (3.5) arise in most models of quark confinement, but with different values of the slopes.

In Fig. 1 we plot the Regge trajectories in the (J, M^2) plane for mesons with natural ($P = (-1)^J$) and unnatural ($P = (-1)^{J-1}$) parity. The Regge trajectories in the (n_r, M^2) plane are presented in Fig. 2. The masses calculated in our model are shown by diamonds. Available experimental data are given by dots (error bars are very small and therefore omitted) with corresponding meson names. Straight lines were obtained by a χ^2 fit of the calculated values. The fitted slopes and intercepts of the Regge trajectories are given in Ref. [4]. We see that the calculated heavy-light meson masses fit nicely to the linear trajectories in both planes. These trajectories are almost parallel and equidistant.

Experimentally complete sets of $1P$-wave meson candidates are known in the charm sector. In the bottom sector masses of only narrow states originating from the $j = 3/2$ heavy quark spin

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Table 1: Masses of charmed ($q = u,d$) and charmed-strange mesons (in MeV).

State		Theory	Experiment [1]		Theory	Experiment [1]	
$n^{2S+1}L_J$	J^P	$c\bar{q}$	meson	mass	$c\bar{s}$	meson	mass
1^1S_0	0^-	1871	D	1869.62(20)	1969	D_s	1968.49(34)
1^3S_1	1^-	2010	$D^*(2010)$	2010.27(17)	2111	D_s^*	2112.3(5)
1^3P_0	0^+	2406	$D_0^*(2400)$	$2403(40)(^{\pm})$ $2318(29)(^{0})$	2509	$D_{s0}^*(2317)$	2317.8(6)
$1P_1$	1^+	2469	$D_1(2430)$	2427(40)	2574	$D_{s1}(2460)$	2459.6(6)
$1P_1$	1^+	2426	$D_1(2420)$	2423.4(3.1)	2536	$D_{s1}(2536)$	2535.35(60)
1^3P_2	2^+	2460	$D_2^*(2460)$	2462.6(0.6)	2571	$D_{s2}(2573)$	2572.6(9)
2^1S_0	0^-	2581	$D(2550)$	2539(8)	2688		
2^3S_1	1^-	2632	$D^*(2600)$ $D^*(2640)$	2612(6) 2637(6)	2731	$D_{s1}(2710)$	$2710(^{+12}_{-7})$
1^3D_1	1^-	2788	$D^*(2760)$	2761(5)	2913		
$1D_2$	2^-	2850			2961		
$1D_2$	2^-	2806			2931		
1^3D_3	3^-	2863			2971	$D_{sJ}^*(2860)$	$2862(^{+6}_{-3})$
2^3P_0	0^+	2919			3054		
$2P_1$	1^+	3021			3154		
$2P_1$	1^+	2932			3067	$D_{sJ}(3040)$	$3044(^{+30}_{-9})$
2^3P_2	2^+	3012			3142		
3^1S_0	0^-	3062			3219		
3^3S_1	1^-	3096			3242		
1^3F_2	2^+	3090			3230		
$1F_3$	3^+	3145			3266		
$1F_3$	3^+	3129			3254		
1^3F_4	4^+	3187			3300		
2^3D_1	1^-	3228			3383		
$2D_2$	2^-	3307			3456		
$2D_2$	2^-	3259			3403		
2^3D_3	3^-	3335			3469		
3^3P_0	0^+	3346			3513		
$3P_1$	1^+	3461			3618		
$3P_1$	1^+	3365			3519		
3^3P_2	2^+	3407			3580		
1^3G_3	3^-	3352			3508		
$1G_4$	4^-	3415			3554		
$1G_4$	4^-	3403			3546		
1^3G_5	5^-	3473			3595		
4^1S_0	0^-	3452			3652		
4^3S_1	1^-	3482			3669		
5^1S_0	0^-	3793			4033		
5^3S_1	1^-	3822			4048		

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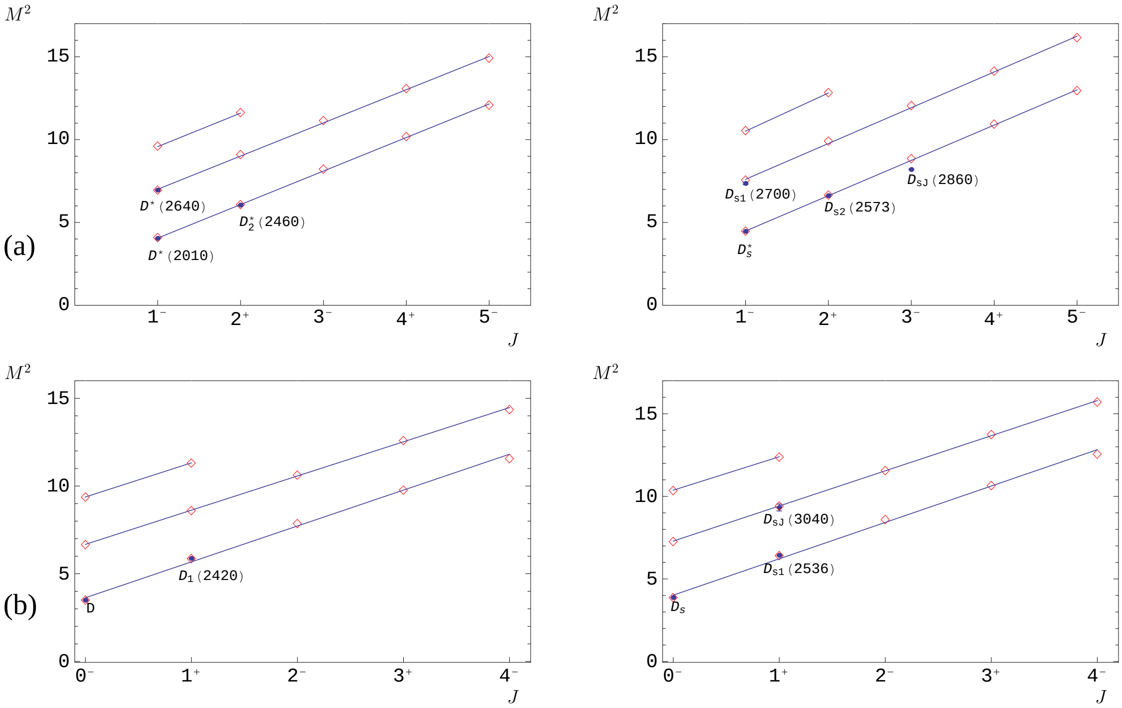

Figure 1: Parent and daughter (J,M^2) Regge trajectories for charmed and charmed-strange mesons with natural (a) and unnatural (b) parity. Diamonds are predicted masses. Available experimental data are given by dots with particle names. M^2 is in GeV^2.

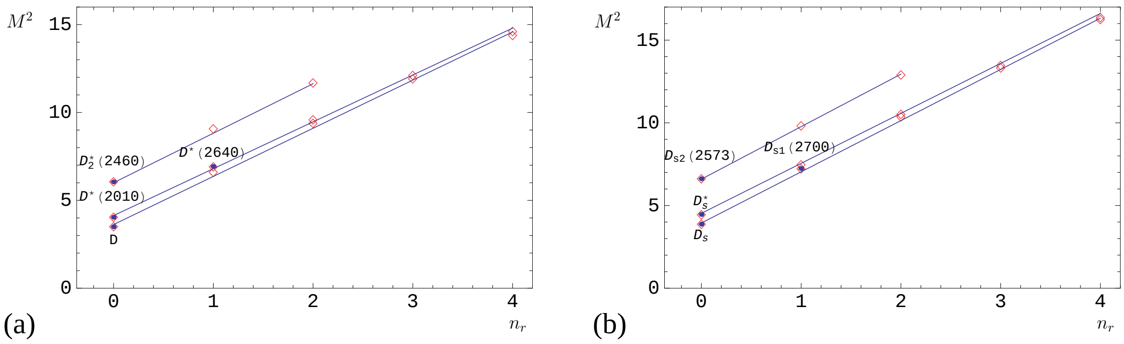

Figure 2: The (n_r,M^2) Regge trajectories for pseudoscalar, vector and tensor charmed (a) and charmed-strange (b) mesons (from bottom to top).

multiplet are known reliably. There are some indications of the broad $j = 1/2$ states both of bottom (0^+) and bottom-strange (1^+) mesons, but additional confirmation is needed. We find good agreement of our predictions for $1P$ wave states with available data except for the masses of $D^*_{s0}(2317)$ and $D_{s1}(2460)$ mesons. These two charmed-strange meson states have anomalously low masses which are even lower than the experimentally observed masses of the corresponding charmed $D^*_0(2400)$ and $D_1(2427)$ mesons. Our model predictions for the masses of the $1P$-wave 0^+ and 1^+ states are almost 200 MeV and 110 MeV higher than the measured masses of $D^*_{s0}(2317)$ and $D_{s1}(2460)$ mesons. Such phenomenon is very hard to understand within the quark-antiquark picture for these states. Most of the explanations available in the literature are based on some very

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specific fine tuning of the model parameters. The influence of such tuning on the spectroscopy of other mesons, which are well described in the framework of the conventional approach, is not well understood. It is probable that these mesons could have an exotic nature and the genuine quark-antiquark P-wave charmed-strange 0^+ and 1^+ states have higher masses above the DK and D^*K thresholds and are, therefore, broad. We find that the $D^*_{s0}(2317)$ and $D_{s1}(2460)$ mesons do not lie on the corresponding Regge trajectories. This may be an additional indication of their anomalous nature. All other experimentally observed $1P$-wave states match well their trajectories.

Our model suggests that $D_{s1}(2700)$ and $D^*(2640)$ mesons are the first radial excitations (2^3S_1) of the vector charmed-strange and charmed mesons. Figures 1 and 2 show that they lie on the corresponding Regge trajectories both in the (J,M^2) and (n_r,M^2) planes.

Recent experimental observation that $D^*_{sJ}(2860)$ decays to both DK and D^*K indicates that this state should have natural parity. In our model natural parity states 1^- (1^3D_1) and 3^- (1^3D_3) have masses which exceed the experimental value by about 50 and 100 MeV, respectively. It was argued that from the point of view of decay rates the 3^- assignment is favored. However the measurement of the branching ratios of the $D^*_{sJ}(2860)$ decay into D^*K to the branching ratio of the decay into DK differs from the theoretical expectations by three standard deviations. From Fig. 1 we see that this state does not fit well to the corresponding Regge trajectory.

On the other hand, the state $D_{sJ}(3040)$, recently observed by BaBar in the D^*K mass spectrum, has a mass coinciding within errors with the mass of the 1^+ ($2P_1$) state predicted by our model (see Table 1). This state nicely fits to the daughter Regge trajectory in Fig. 1

4. Heavy baryons

The heavy-quark–light-diquark picture of heavy baryons reduces the calculations to two steps. The first step is the calculation of masses, wave functions and form factors of the diquarks, composed from two relativistic light quarks. We follow the same approach which was previously used for heavy-light mesons (3.1). Next, at the second step, a heavy baryon is treated as a bound system of the relativistic light diquark and heavy quark. It is important to emphasize that we explicitly take into account the diquark structure through the diquark-gluon vertex expressed in terms of the diquark wave function. Then the vertex of the diquark-gluon interaction in (2.7) is given by

$$\langle d(P)|J_\mu(0)|d(Q)\rangle = \int \frac{d^3p\, d^3q}{(2\pi)^6}\bar{\Psi}^d_P(\mathbf{p})\Gamma_\mu(\mathbf{p},\mathbf{q})\Psi^d_Q(\mathbf{q}). \tag{4.1}$$

It leads to the emergence of the diquark-gluon form factor $F(r)$ which with high accuracy can be approximated by the expression [5]

$$F(r) = 1 - e^{-\xi r - \zeta r^2}. \tag{4.2}$$

The values of the diquark masses and parameters ξ and ζ for the ground states of the scalar $[q,q']$ and axial vector $\{q,q'\}$ light diquarks are given in Ref. [6].

We solve numerically the quasipotential equation with the quasipotential which nonperturbatively accounts for the relativistic dynamics both of the light diquark d and heavy quark Q. The calculated values of the ground state and excited baryon masses are given in Table 2 for Λ_Q and Ξ_Q baryons with the scalar diquark (for masses of other baryons see Ref. [6]) in comparison with

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available experimental data [1]. In the first two columns we give the baryon quantum numbers (J^P) and the state of the heavy-quark–light-diquark (Qd) bound system (in usual notations $(n_r+1)L$), while in the remaining columns our predictions for the masses and experimental data are shown.

It is important to note that in the adopted quark-diquark picture of heavy baryons we consider solely the orbital and radial excitations of the bound heavy quark and light diquark, while light diquarks are taken in the ground (scalar or axial-vector) state. As a result, we get significantly less excited states than in the genuine three-quark picture of a baryon. As it is seen from Table 2, such an approach is supported by available experimental data, which are nicely accommodated in the quark-diquark picture.

In Fig. 3 we plot, as an example, the Regge trajectories in the (J,M^2) plane for charmed and bottom baryons with natural ($P=(-1)^{J-1/2}$) and unnatural ($P=(-1)^{J+1/2}$) parities. Other Regge trajectories including the ones in the (n_r,M^2) plane are presented in Ref. [6]. Again we see that the calculated heavy baryon masses fit nicely to the linear trajectories in both planes. These trajectories are almost parallel and equidistant.

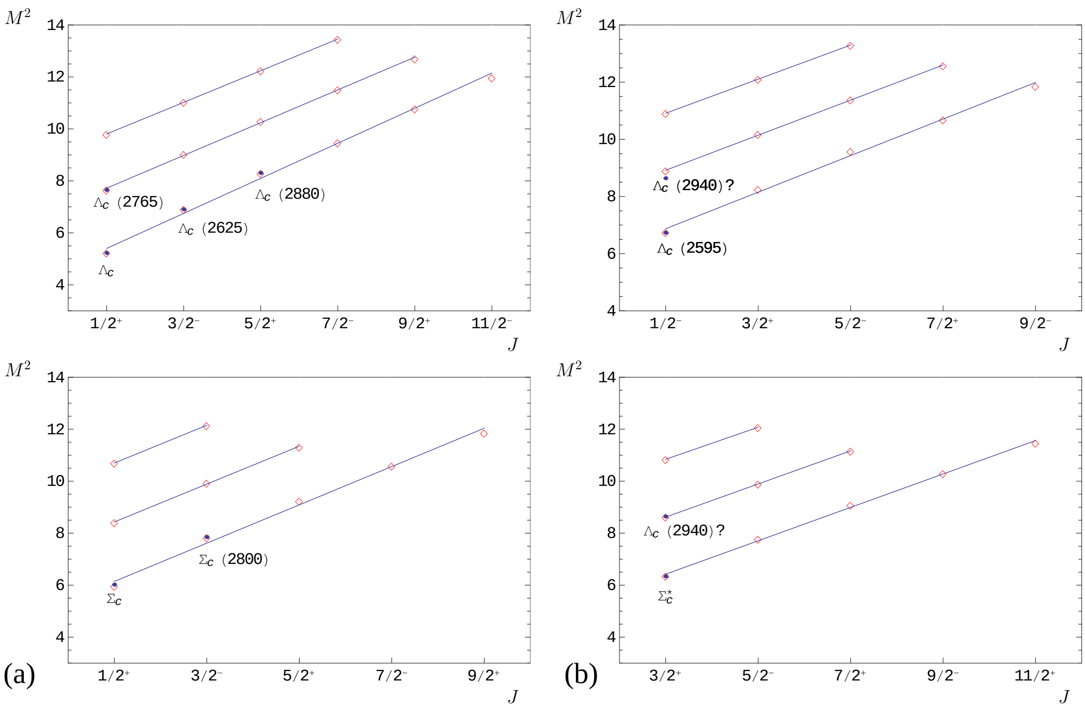

Figure 3: Parent and daughter (J,M^2) Regge trajectories for the Λ_c and Σ_c baryons with natural (a) and unnatural (b) parities. Diamonds are predicted masses. Available experimental data are given by dots with particle names; M^2 is in GeV^2.

The obtained results allow us to determine the possible quantum numbers of the observed heavy baryons and prescribe them to a particular Regge trajectory. In the (J,M^2) plane there are three trajectories for which three experimental candidates are available (parent trajectories for the $\Lambda_c\left(\frac{1}{2}^+\right)$, for the $\Xi_c\left(\frac{1}{2}^+\right)$ and for the $\Xi_c^*\left(\frac{3}{2}^+\right)$) and two trajectories with two experimental candidates (parent trajectories for the $\Sigma_c\left(\frac{1}{2}^+\right)$ and for the $\Xi_c\left(\frac{1}{2}^-\right)$). On the other hand, in the (n_r,M^2) plane there are three trajectories with two experimental candidates (the $\Lambda_c\left(\frac{1}{2}^+\right)$ and the $\Lambda_c\left(\frac{1}{2}^-\right)$

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Table 2: Masses of the Λ_Q and Ξ_Q ($Q=c,b$) heavy baryons with the scalar diquark (in MeV).

	Qd	Λ_c		Λ_b		Ξ_c		Ξ_b	
J^P	state	M	$M^{\rm exp}$ [1]	M	$M^{\rm exp}$ [1]	M	$M^{\rm exp}$ [1]	M	$M^{\rm exp}$ [1]
$\frac{1}{2}^+$	$1S$	2286	2286.46(14)	5620	5620.2(1.6)	2476	$2470.88(^{34}_{80})$	5803	5790.5(2.7)
$\frac{1}{2}^+$	$2S$	2769	2766.6(2.4)?	6089		2959		6266	
$\frac{1}{2}^+$	$3S$	3130		6455		3323		6601	
$\frac{1}{2}^+$	$4S$	3437		6756		3632		6913	
$\frac{1}{2}^+$	$5S$	3715		7015		3909		7165	
$\frac{1}{2}^+$	$6S$	3973		7256		4166		7415	
$\frac{1}{2}^-$	$1P$	2598	2595.4(6)	5930		2792	2791.8(3.3)	6120	
$\frac{1}{2}^-$	$2P$	2983	$2939.3(^{1.4}_{1.5})$?	6326		3179		6496	
$\frac{1}{2}^-$	$3P$	3303		6645		3500		6805	
$\frac{1}{2}^-$	$4P$	3588		6917		3785		7068	
$\frac{1}{2}^-$	$5P$	3852		7157		4048		7302	
$\frac{3}{2}^-$	$1P$	2627	2628.1(6)	5942		2819	2819.6(1.2)	6130	
$\frac{3}{2}^-$	$2P$	3005		6333		3201		6502	
$\frac{3}{2}^-$	$3P$	3322		6651		3519		6810	
$\frac{3}{2}^-$	$4P$	3606		6922		3804		7073	
$\frac{3}{2}^-$	$5P$	3869		7171		4066		7306	
$\frac{3}{2}^+$	$1D$	2874		6190		3059	3054.2(1.3)	6366	
$\frac{3}{2}^+$	$2D$	3189		6526		3388		6690	
$\frac{3}{2}^+$	$3D$	3480		6811		3678		6966	
$\frac{3}{2}^+$	$4D$	3747		7060		3945		7208	
$\frac{5}{2}^+$	$1D$	2880	2881.53(35)	6196		3076	3079.9(1.4)	6373	
$\frac{5}{2}^+$	$2D$	3209		6531		3407		6696	
$\frac{5}{2}^+$	$3D$	3500		6814		3699		6970	
$\frac{5}{2}^+$	$4D$	3767		7063		3965		7212	
$\frac{5}{2}^-$	$1F$	3097		6408		3278		6577	
$\frac{5}{2}^-$	$2F$	3375		6705		3575		6863	
$\frac{5}{2}^-$	$3F$	3646		6964		3845		7114	
$\frac{5}{2}^-$	$4F$	3900		7196		4098		7339	
$\frac{7}{2}^-$	$1F$	3078		6411		3292		6581	
$\frac{7}{2}^-$	$2F$	3393		6708		3592		6867	
$\frac{7}{2}^-$	$3F$	3667		6966		3865		7117	
$\frac{7}{2}^-$	$4F$	3922		7197		4120		7342	
$\frac{7}{2}^+$	$1G$	3270		6598		3469		6760	
$\frac{7}{2}^+$	$2G$	3546		6867		3745		7020	
$\frac{9}{2}^+$	$1G$	3284		6599		3483		6762	
$\frac{9}{2}^+$	$2G$	3564		6868		3763		7032	
$\frac{9}{2}^-$	$1H$	3444		6767		3643		6933	
$\frac{11}{2}^-$	$1H$	3460		6766		3658		6934	

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and the $\Xi_c\left(\frac{1}{2}^+\right)$. All experimental points fit well to the corresponding Regge trajectories obtained in our model.

From calculated masses (Table 2) and plotted Regge trajectories (Fig. 3) we see that the $\Lambda_c(2765)$ (or $\Sigma_c(2765)$), [1] if it is indeed the Λ_c state, can be interpreted in our model as the first radial ($2S$) excitation of the Λ_c. If instead it is the Σ_c state, then it can be identified as its first orbital excitation ($1P$) with $J=\frac{3}{2}^-$. The $\Lambda_c(2880)$ baryon corresponds to the second orbital excitation ($2D$) with $J=\frac{5}{2}^+$, fitting nicely the parent Λ_c Regge trajectory in the (J,M^2) plane (see Fig. 3). Such prescription is in accord with the experimental evidence coming from the $\Sigma_c(2455)\pi$ decay angular distribution [1]. The other charmed baryon, denoted as $\Lambda_c(2940)$, probably has $I=0$, since it was discovered in the pD^0 mass spectrum and not observed in the pD^+ channel, but $I=1$ is not ruled out [1]. If it is really the Λ_c, state then it could be an orbitally and radially excited ($2P$) state with $J=\frac{1}{2}^-$, whose mass is predicted to be about 40 MeV heavier (see Fig. 3). A better agreement with experiment (within few MeV) is achieved, if the $\Lambda_c(2940)$ is interpreted as the first radial excitation ($2S$) of the Σ_c with $J=\frac{3}{2}^+$. The $\Sigma_c(2800)$ can be identified with one of the first orbital ($1P$) excitations of the Σ_c with $J=\frac{1}{2}^-$ or $\frac{3}{2}^-$ which have very close masses compatible with experimental value within errors.

From the results for masses and the Regge trajectories of the Ξ_Q baryons both with the scalar and axial vector diquarks we see that the $\Xi_c(2790)$ and $\Xi_c(2815)$ can be assigned to the first orbital ($1P$) excitations of the Ξ_c containing a scalar diquark with $J=\frac{1}{2}^-$ and $J=\frac{3}{2}^-$, respectively. On the other hand, the charmed baryon $\Xi_c(2930)$ can be considered as either the $J=\frac{1}{2}^-$, $J=\frac{3}{2}^-$ or $J=\frac{5}{2}^-$ state (all these states are predicted to have close masses) corresponding to the first orbital ($1P$) excitations of the Ξ_c' with an axial vector diquark. While the $\Xi_c(2980)$ can be viewed as the first radial ($2S$) excitation with $J=\frac{1}{2}^+$ of the Ξ_c', the $\Xi_c(3055)$ and $\Xi_c(3080)$ baryons can be interpreted as a second orbital ($2D$) excitations of the Ξ_c containing a scalar diquark with $J=\frac{3}{2}^+$ and $J=\frac{5}{2}^+$, and the $\Xi_c(3123)$ can be viewed as the corresponding ($2D$) excitation of the Ξ_c' with $J=\frac{7}{2}^+$.

For the Ω_c baryons as well as for all bottom baryons only masses of ground states are known [1], most of which were measured recently. Our original predictions for the ground states [5] of these baryons are very close to the values presented in [6] and agree well with measurements [1].

It is interesting to compare the values of the slopes of the Regge trajectories for heavy baryons [6], heavy-light [4] and light mesons [3]. Comparing corresponding values, we find that for the same flavour of the heavy quark the heavy baryon slopes have higher values than the heavy-light meson ones. It is also worth remarking that the ratios of the heavy baryon to heavy-light meson slopes ($\alpha_{Qqq}/\alpha_{Q\bar{q}}$ and $\beta_{Qqq}/\beta_{Q\bar{q}}$) have very close values, which are about 1.4, both in the (J,M^2) and in (n_r,M^2) planes. Note that light baryons and light mesons have almost equal values of the Regge trajectory slopes.

[1] It is important to note that the J^P quantum numbers for most excited heavy baryons have not been determined experimentally, but are assigned by PDG on the basis of quark model predictions. For some excited charm baryons such as the $\Lambda_c(2765)$, $\Lambda_c(2880)$ and $\Lambda_c(2940)$ it is even not known if they are excitations of the Λ_c or Σ_c.

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5. Conclusions

In this talk the spectroscopy of charmed and bottom mesons and baryons was considered in the framework of the quark-diquark picture in the relativistic quark model. The heavy baryon was treated as a heavy-quark–light-diquark bound system in which excitations occur only between a heavy quark and a light diquark. The light diquarks were considered only in the ground (either scalar or axial vector) state. The diquark internal structure was taken into account by including form factors of the diquark-gluon interaction in terms of the diquark wave functions. The dynamics of light and heavy quarks and diquark was treated completely relativistically without applying of either the nonrelativistic v/c or heavy quark $1/m_Q$ expansions. Such nonperturbative approach is especially important for the highly excited charmed hadron states, where the heavy quark expansion is not adequate.

We presented the calculated masses of ground state, orbitally and radially excited heavy mesons and baryons up to rather high excitations. This allowed us to construct the Regge trajectories both in the (J,M^2) and (n_r,M^2) planes. It was found that they are almost linear, parallel and equidistant. Most of the available experimental data nicely fit to them. In the meson sector, the anomalously light $D^*_{s0}(2317)$, $D_{s1}(2460)$ and $D^*_{sJ}(2860)$ mesons, which masses are 100-200 MeV lower than various model predictions, are exceptions. The masses of the charmed-strange $D^*_{s0}(2317)$, $D_{s1}(2460)$ mesons either almost coincide or are even lower than the masses of the partner charmed $D^*_0(2400)$ and $D_1(2427)$ mesons. These states thus could have an exotic origin [7]. It will be very important to find the bottom counterparts of these states in order to reveal their nature. On the other hand, in the baryon sector the assignment of the experimentally observed heavy baryons to the particular Regge trajectories allowed us to suggest the quantum numbers of the excited heavy baryons. It was found that all currently available experimental data can be well described in the relativistic quark-diquark picture, which predicts significantly less states than the genuine three-body picture.

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