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Baryon Properties in the Relativistic Quark Model[1]

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Abstract—Properties of heavy and strange baryons are investigated in the framework of the relativistic quark-diquark picture. It is based on the relativistic quark model of hadrons, which was previously successfully applied for the calculation of meson properties. It is assumed that two quarks in a baryon form a diquark and baryon is considered as the bound quark-diquark system. The relativistic effects and diquark internal structure are consistently taken into account. Calculations are performed up to rather high orbital and radial excitations of heavy and strange baryons. On this basis the Regge trajectories are constructed. The rates of semileptonic decays of heavy baryons are calculated. The obtained results agree well with available experimental data.

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The convincing evidence of the existence of diquark correlations in hadrons has been collected. Recently several charged charmonium- and bottomonium-like states were discovered. They should be inevitably multiquark, at least four quark states. One of the most successful pictures of such tetraquark states is the diquark-antidiquark model [1]. In the light meson sector it has been argued for a long time that mesons forming the inverted lightest scalar nonet can be well described as tetraquarks treated as diquark-antidiquark bound states [2]. In the baryon sector it is well known that the number of observed excited states both in the light and heavy sectors is considerably lower than the number of excited states predicted in the three-quark picture [3]. The introduction of diquarks significantly reduces this number since in such a picture some of degrees of freedom are frozen and thus the number of possible excitations is substantially smaller.

Here we study spectroscopy and weak decays of heavy and strange baryons in the relativistic quark-diquark picture in the framework of the quasipotential approach. The interaction of two quarks in a diquark and the quark-diquark interaction in a baryon are described by the diquark wave function Ψ_d of the bound quark-quark state and by the baryon wave function Ψ_B of the bound quark-diquark state respectively, which satisfy the relativistic quasipotential equation of the Schrödinger type [4]

$$\left(\frac{b^2(M)}{2\mu_R}-\frac{\mathbf{p}^2}{2\mu_R}\right)\Psi_{d,B}(\mathbf{p}) = \int\frac{d^3q}{(2\pi)^3}V(\mathbf{p},\mathbf{q};M)\Psi_{d,B}(\mathbf{q}), \tag{1}$$

where the relativistic reduced mass and the center-of-mass system relative momentum squared on mass shell are

$$\mu_R=\frac{M^4-(m_1^2-m_2^2)^2}{4M^3},$$
$$b^2(M)=\frac{[M^2-(m_1+m_2)^2][M^2-(m_1-m_2)^2]}{4M^2},$$

and M is the bound state mass (diquark or baryon), $m_{1,2}$ are the masses of quarks (q_1 and q_2) which form the diquark or of the diquark (d) and quark (q) which form the baryon (B), and $\mathbf{p}$ is their relative momentum.

The kernel $V(\mathbf{p},\mathbf{q};M)$ in Eq. (1) is the quasipotential operator of the quark-quark or quark-diquark interaction which is constructed with the help of the off-mass-shell scattering amplitude, projected onto the positive energy states. We assume that the effective interaction is the sum of the usual one-gluon exchange term and the mixture of long-range vector and scalar linear confining potentials, where the vector confining potential contains the Pauli term. The resulting quasipotentials are given by the following expressions.

(a) Quark-quark (qq) interaction in the diquark

$$V(\mathbf{p},\mathbf{q};M)=\bar{u}_1(p)\bar{u}_2(-p)\frac{1}{2}\left[\frac{4}{3}\alpha_sD_{\mu\nu}(\mathbf{k})\gamma_1^\mu\gamma_2^\nu + V_{\rm conf}^V(\mathbf{k})\Gamma_1^\mu(\mathbf{k})\Gamma_{2;\mu}(-\mathbf{k})+V_{\rm conf}^S(\mathbf{k})\right]u_1(q)u_2(-q). \tag{2}$$

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(b) Quark-diquark (qd) interaction in the baryon

$$V(\mathbf{p},\mathbf{q};M) = \frac{\langle d(P)|J_\mu|d(Q)\rangle}{2\sqrt{E_d(p)E_d(q)}}\bar{u}_q(p)\frac{4}{3}\alpha_s D_{\mu\nu}(\mathbf{k})\gamma^\nu u_q(q)$$
$$+ \psi_d^*(P)\bar{u}_q(p)J_{d;\mu}\Gamma_q^\mu(\mathbf{k})V_{\rm conf}^V(\mathbf{k})u_q(q)\psi_d(Q) \quad (3)$$
$$+ \psi_d^*(P)\bar{u}_q(p)V_{\rm conf}^S(\mathbf{k})u_q(q)\psi_d(Q),$$

where α_s is the QCD coupling constant, $\langle d(P)|J_\mu|d(Q)\rangle$ is the vertex of the diquark-gluon interaction which takes into account the diquark internal structure and $J_{d;\mu}$ is the effective long-range vector vertex of the diquark. The diquark momenta are $P = (E_d(p),-\mathbf{p})$, $Q = (E_d(q),-\mathbf{q})$ with $E_d(p) = \sqrt{\mathbf{p}^2 + M_d^2}$. $D_{\mu\nu}$ is the gluon propagator in the Coulomb gauge, $\mathbf{k} = \mathbf{p} - \mathbf{q}$; γ_μ and $u(p)$ are the Dirac matrices and spinors, while $\psi_d(P)$ is the diquark wave function. The factor 1/2 in the quark-quark interaction accounts for the difference of the colour factor compared to the quark-antiquark case.

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

$$\Gamma_\mu(\mathbf{k}) = \gamma_\mu + \frac{i\kappa}{2m}\sigma_{\mu\nu}\tilde{k}^\nu, \quad \tilde{k} = (0,\mathbf{k}), \quad (4)$$

where κ is the anomalous chromomagnetic moment of quarks.

In the nonrelativistic limit the vector and scalar confining potentials reduce to

$$V_{\rm conf}^V(r) = (1-\varepsilon)(Ar+B), \quad V_{\rm conf}^S(r) = \varepsilon(Ar+B), \quad (5)$$

where ε is the mixing coefficient, and the usual Cornell-like potential is reproduced

$$V(r) = -\frac{4}{3}\frac{\alpha_s}{r} + Ar + B. \quad (6)$$

Here we use the QCD coupling constant with freezing

$$\alpha_s(\mu^2) = \frac{4\pi}{\beta_0 \ln\frac{\mu^2 + M_B^2}{\Lambda^2}}, \quad \beta_0 = 11 - \frac{2}{3}n_f,$$
$$\mu = \frac{2m_1m_2}{m_1+m_2}, \quad (7)$$

with the background mass $M_B = 2.24\sqrt{A} = 0.95$ GeV and $\Lambda = 413$ MeV [2].

All parameters of the model such as quark masses, parameters of the linear confining potential A and B, the mixing coefficient ε and anomalous chromomagnetic quark moment κ were fixed previously from calculations of meson and baryon properties [4, 5]. The constituent quark masses $m_u = m_d = 0.33$ GeV, $m_s = 0.5$ GeV and the parameters of the linear potential $A = 0.18$ GeV^2 and $B = -0.3$ GeV have the usual values of quark models. The value of the mixing coefficient of vector and scalar confining potentials $\varepsilon = -1$ has been determined from the consideration of the heavy quark expansion for the semileptonic heavy meson decays and charmonium radiative decays [4]. While the universal Pauli interaction constant $\kappa = -1$ has been fixed from the analysis of the fine splitting of heavy quarkonia ^3P_J-states [4]. Note that the long-range chromomagnetic contribution to the potential, which is proportional to $(1+\kappa)$, vanishes for the chosen value of $\kappa = -1$.

Table 1. Masses M and form factor parameters of diquarks

Quark content	I	State nl_j	M, MeV	ξ, GeV	ζ, GeV^2
ud	0	$1s_0$	710	1.09	0.185
	1	$1s_1$	909	1.185	0.365
	0	$1p_0$	1321	1.395	0.148
	0	$1p_1$	1397	1.452	0.195
	0	$1p_2$	1475	1.595	0.173
	1	$1p_1$	1392	1.451	0.194
	0	$2s_0$	1513	1.01	0.055
	1	$2s_1$	1630	1.05	0.151
ss	0	$1s_1$	1203	1.13	0.280
	0	$1p_1$	1608	1.03	0.208
	0	$2s_1$	1817	0.805	0.235

First we calculate masses and form factors of the diquarks. The obtained masses of the ground and excited states of diquarks are presented in Table 1. In this table we also give the values of the parameters ξ and ζ. They enter the vertex $\langle d(P)|J_\mu|d(Q)\rangle$ of the diquark-gluon interaction (3) which is parameterized by the form factor $F(r) = 1 - e^{-\xi r - \zeta r^2}$, that takes the internal diquark structure into account.

Next we calculate the masses of heavy and strange baryons as the bound states of a quark and diquark. The results of calculations can be found in Refs. [5, 6]. Here, as an example, we give in Tables 2, 3 the comparison of our predictions for the mass spectra of the Λ_c and Λ with other theoretical estimates and available experimental data. From these tables we see that our diquark model predicts appreciably less states than the three-quark approaches. The differences become apparent with the growth of the orbital and radial excitations in the baryon. Our model reproduces correctly all masses of the well established 4- and 3-star resonances and most of the 2- and 1-star states. Let us emphasize that the experimental mass of the $\frac{1}{2}^-$ 4-star $\Lambda(1405)$ is naturally reproduced if this state is considered as the first orbital excitation in the strange quark-

Table 2. Comparison of theoretical predictions for masses of the Λ_c baryons (in MeV)

J^P	Experiment [7]			Theory			
	state	status	mass	our	[8]	[3]	[9]
$\frac{1}{2}^+$	Λ_c $\Lambda_c(2765)$	*** *	2286.46(14) 2766.6(2.4)	2286 2769 3130 3437	2286 2766 3112 3397	2286 2791 3154	2265 2775 3170
$\frac{1}{2}^-$	$\Lambda_c(2595)$ $\Lambda_c(2940)$	*** ***	2592.25(28) $2939.3(^{1.4}_{1.3})$	2598 2983 3303	2591 2989 3296	2625	2630 2780 2830
$\frac{3}{2}^-$	$\Lambda_c(2625)$	***	2628.1(6)	2627 3005 3322	2629 3000 3301	2636	2640 2840 2885
$\frac{3}{2}^+$				2874 3189	2857 3188	2887 3120	2910 3035
$\frac{5}{2}^+$	$\Lambda_c(2880)$	***	2881.53(35)	2880 3209	2879 3198	2887 3125	2910 3140
$\frac{5}{2}^-$				3097	3075	2872	2900
$\frac{7}{2}^-$				3078	3092		3125
$\frac{7}{2}^+$				3270	3267		3175
$\frac{9}{2}^+$				3284	3280		
$\frac{9}{2}^-$				3444			
$\frac{11}{2}^-$				3460			

light scalar diquark picture of Λ baryons. The rather low mass of this state represents difficulties for most of the three-quark models [9–11], which predict its mass about 100 MeV higher than experimental value.

We calculated masses of both orbitally and radially excited heavy and strange baryons up to rather high excitation numbers. This makes it possible to construct the heavy baryon Regge trajectories both in the (J, M^2) and in the (n_r, M^2) planes: (a) The (J, M^2) Regge trajectory: $J = \alpha M^2 + \alpha_0$; (b) The (n_r, M^2) Regge trajectory: $n_r = \beta M^2 + \beta_0$, where α, β are the slopes and α_0, β_0 are intercepts. In Fig. 1 we plot the Regge trajectories in the (J, M^2) plane for charmed and strange baryons. Straight lines were obtained by the χ^2 fit of calculated values. The fitted slopes and intercepts of the Regge trajectories can be found in Refs. [5, 6]. We see that the calculated heavy and strange baryon masses lie on the linear trajectories, which are parallel and equidistant.

Next we consider weak decays of baryons. The simplest case is the semileptonic decay of the bottom baryon to the charmed baryon, since in this case the model independent considerations based on heavy quark effective theory (HQET) can be applied. To calculate the heavy baryon decay rate it is necessary to determine the corresponding matrix element of the weak current between baryon states, which in the quasipotential approach is given by

$$\left\langle B_{Q'}(p_{Q'})\middle|J^W_\mu\middle|B_Q(p_Q)\right\rangle = \int\frac{d^3p\,d^3q}{(2\pi)^6}\bar{\Psi}_{B_{Q'}\mathbf{p}_{Q'}}(\mathbf{p})\Gamma_\mu(\mathbf{p},\mathbf{q})\Psi_{B_Q\mathbf{p}_Q}(\mathbf{q}), \tag{8}$$

where $\Gamma_\mu(\mathbf{p},\mathbf{q})$ is the two-particle vertex function and $\Psi_{B\mathbf{p}_Q}$ are the baryon ($B = B_Q, B_{Q'}$) wave functions pro-

Table 3. Comparison of theoretical predictions and experimental data for the masses of the Λ states (in MeV)

J^P	Experiment [7]			Theory					
	state	status	mass	our	[9]	[10]	[11]	[12]	[13]
$\frac{1}{2}^+$	Λ	****	1115.683(6)	1115	1115	1108	1136	1116	1149(18)
	Λ(1600)	***	1560–1600	1615	1680	1677	1625	1518	1807(94)
	Λ(1710)	*	1713(13)						
	Λ(1810)	***	1750–1810	1901	1830	1747	1799	1666	2112(54)
				1972	1910	1898		1955	2137(69)
				1986	2010	2077		1960	
				2042	2105	2099			
				2099	2120	2132			
$\frac{3}{2}^+$	Λ(1890)	****	1850–1890	1854	1900	1823		1896	1991(103)
				1976	1960	1952			2058(139)
				2130	1995	2045			2481(111)
				2184	2050	2087			
				2202	2080	2133			
$\frac{5}{2}^+$	Λ(1820)	****	1815–1820	1825	1890	1834		1896	
	Λ(2110)	***	2090–2110	2098	2035	1999			
				2221	2115	2078			
				2255	2115	2127			
				2258	2180	2150			
$\frac{7}{2}^+$	Λ(2020)	*	≈ 2020	2251	2120	2130			
				2471		2331			
$\frac{9}{2}^+$	Λ(2350)	***	2340–2350	2360		2340			
$\frac{1}{2}^-$	Λ(1405)	****	$1405.1(^{1.3}_{1.0})$	1406	1550	1524	1556	1431	1416(81)
	Λ(1670)	****	1660–1670	1667	1615	1630	1682	1443	1546(110)
	Λ(1800)	***	1720–1800	1733	1675	1816	1778	1650	1713(116)
				1927	2015	2011		1732	2075(249)
				2197	2095	2076		1785	
				2218	2160	2117		1854	
$\frac{3}{2}^-$	Λ(1520)	****	1519.5(1.0)	1549	1545	1508	1556	1431	1751(40)
	Λ(1690)	****	1685–1690	1693	1645	1662	1682	1443	2203(106)
				1812	1770	1775		1650	2381(87)
	Λ(2050)	*	2056(22)	2035	2030	1987		1732	
				2319	2110	2090		1785	
	Λ(2325)	*	≈ 2325	2322	2185	2147		1854	
				2392	2230	2259		1928	
				2454	2290	2275		1969	
$\frac{5}{2}^-$	Λ(1830)	****	1810–1830	1861	1775	1828	1778	1785	
				2136	2180	2080			
				2350	2250	2179			
$\frac{7}{2}^-$	Λ(2100)	****	2090–2100	2097	2150	2090			
				2583	2230	2227			
$\frac{9}{2}^-$				2665		2370			

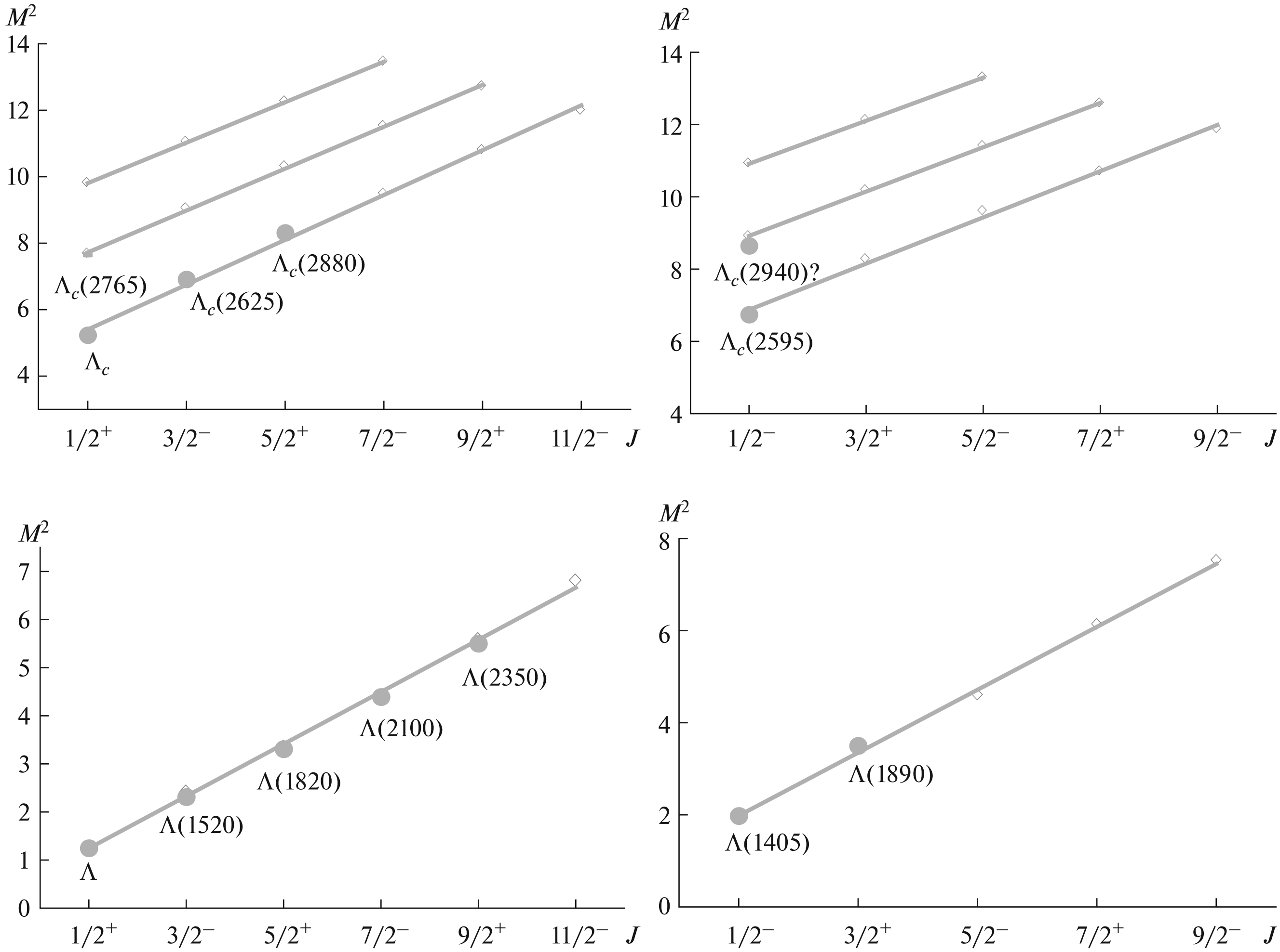

Fig. 1. Parent and daughter (J, M^2) Regge trajectories for the Λ_c and Λ baryons with natural ($P = -1)^{J-1/2}$) (a) and unnatural ($P = (-1)^{J+1/2}$) (b) parities. Diamonds are predicted masses. Available experimental data are given by dots with particle names; M^2 is in GeV^2.

jected onto the positive energy states of quarks and boosted to the moving reference frame with momentum $\mathbf{p}_Q$.

The hadronic matrix elements for the semileptonic decay $\Lambda_Q \to \Lambda_{Q'}$ are parameterized in terms of six invariant form factors:

$$\begin{aligned}&\left\langle \Lambda_{Q'}(v')\middle|V^{\mu}\middle|\Lambda_Q(v)\right\rangle\\ &= \bar{u}_{\Lambda_{Q'}}(v')[F_1(w)\gamma^{\mu} + F_2(w)v^{\mu} + F_3(w)v'^{\mu}]u_{\Lambda_Q}(v),\\ &\left\langle \Lambda_{Q'}(v')\middle|A^{\mu}\middle|\Lambda_Q(v)\right\rangle\\ &= \bar{u}_{\Lambda_{Q'}}(v')[G_1(w)\gamma^{\mu} + G_2(w)v^{\mu} + G_3(w)v'^{\mu}]\gamma_5 u_{\Lambda_Q}(v),\end{aligned} \quad (9)$$

where $u_{\Lambda_Q}(v)$ and $u_{\Lambda_{Q'}}(v')$ are Dirac spinors of the initial and final baryon with four-velocities v and v', respectively; $w = v\cdot v'$. In the heavy quark limit $m_Q \to \infty$ $(Q = b, c)$ the form factors (9) can be expressed through the single Isgur–Wise function $\zeta(w)$

$$\begin{aligned}F_1(w) &= G_1(w) = \zeta(w),\\ F_2(w) &= F_3(w) = G_2(w) = G_3(w) = 0.\end{aligned} \quad (10)$$

At subleading order of the heavy quark expansion the form factors are given by

$$\begin{aligned}F_1(w) &= \zeta(w) + \left(\frac{\bar{\Lambda}}{2m_Q} + \frac{\bar{\Lambda}}{2m_{Q'}}\right)[2\chi(w) + \zeta(w)],\\ \bar{\Lambda} &= M_{\Lambda_Q} - m_Q,\\ G_1(w) &= \zeta(w) + \left(\frac{\bar{\Lambda}}{2m_Q} + \frac{\bar{\Lambda}}{2m_{Q'}}\right)\\ &\quad\times\left[2\chi(w) + \frac{w-1}{w+1}\zeta(w)\right],\\ F_2(w) &= G_2(w) = -\frac{\bar{\Lambda}}{2m_{Q'}}\frac{2}{w+1}\zeta(w),\\ F_3(w) &= -G_3(w) = -\frac{\bar{\Lambda}}{2m_Q}\frac{2}{w+1}\zeta(w).\end{aligned} \quad (11)$$

The resulting expressions for the semileptonic decay $\Lambda_Q \to \Lambda_{Q'}$ form factors up to subleading order in $1/m_Q$ in our model are the following [14]

$$F_1(w) = \zeta(w) + \left(\frac{\bar\Lambda}{2m_Q} + \frac{\bar\Lambda}{2m_{Q'}}\right)[2\chi(w) + \zeta(w)]$$
$$+ 4(1-\varepsilon)(1+\kappa)\left[\frac{\bar\Lambda}{2m_{Q'}}\frac{1}{w-1} - \frac{\bar\Lambda}{2m_Q}(w+1)\right]\chi(w),$$
$$G_1(w) = \zeta(w) + \left(\frac{\bar\Lambda}{2m_Q} + \frac{\bar\Lambda}{2m_{Q'}}\right)\left[2\chi(w) + \frac{w-1}{w+1}\zeta(w)\right]$$
$$- 4(1-\varepsilon)(1+\kappa)\frac{\bar\Lambda}{2m_Q}w\chi(w),$$
$$F_2(w) = -\frac{\bar\Lambda}{2m_{Q'}}\frac{2}{w+1}\zeta(w) - 4(1-\varepsilon)(1+\kappa)$$
$$\times\left[\frac{\bar\Lambda}{2m_{Q'}}\frac{1}{w-1} + \frac{\bar\Lambda}{2m_Q}w\right]\chi(w), \quad (12)$$
$$F_2(w) = -\frac{\bar\Lambda}{2m_{Q'}}\frac{2}{w+1}\zeta(w) - 4(1-\varepsilon)(1+\kappa)$$
$$\times\left[\frac{\bar\Lambda}{2m_{Q'}}\frac{1}{w-1} + \frac{\bar\Lambda}{2m_Q}w\right]\chi(w),$$
$$G_2(w) = -\frac{\bar\Lambda}{2m_{Q'}}\frac{2}{w+1}\zeta(w)$$
$$- 4(1-\varepsilon)(1+\kappa)\frac{\bar\Lambda}{2m_{Q'}}\frac{1}{w-1}\chi(w),$$
$$F_3(w) = -G_3(w) = -\frac{\bar\Lambda}{2m_Q}\frac{2}{w+1}\zeta(w)$$
$$+ 4(1-\varepsilon)(1+\kappa)\frac{\bar\Lambda}{2m_Q}\chi(w),$$

where the leading order Isgur–Wise function of heavy baryons

$$\zeta(w) = \lim_{m_Q\to\infty}\int\frac{d^3p}{(2\pi)^3}\Psi_{\Lambda_{Q'}}\left(\mathbf{p} + 2\epsilon_d(p)\sqrt{\frac{w-1}{w+1}}\mathbf{e}_\Delta\right)\Psi_{\Lambda_Q}(\mathbf{p}),$$
$$\mathbf{e}_\Delta = \Delta/\sqrt{\Delta^2}, \quad (13)$$

$\Delta = M_{\Lambda_{Q'}}\mathbf{v}' - M_{\Lambda_Q}\mathbf{v}$, $\epsilon_d(p) = \sqrt{\mathbf{p}^2 + M_d^2}$ and the subleading function

$$\chi(w) = -\frac{w-1}{w+1}\lim_{m_Q\to\infty}\int\frac{d^3p}{(2\pi)^3}\frac{\bar\Lambda - \epsilon_d(p)}{2\bar\Lambda}$$
$$\times\Psi_{\Lambda_{Q'}}\left(\mathbf{p} + 2\epsilon_d(p)\sqrt{\frac{w-1}{w+1}}\mathbf{e}_\Delta\right)\Psi_{\Lambda_Q}(\mathbf{p}). \quad (14)$$

These functions were calculated using our model wave functions of the Λ_b and Λ_c baryons obtained in their mass spectrum calculations.

For $(1-\varepsilon)(1+\kappa) = 0$ the HQET results (11) are reproduced. This can be achieved either setting $\varepsilon = 1$ (pure scalar confinement) or $\kappa = -1$. In our model we need vector confining contribution and therefore use the latter option. This consideration gives us additional, based on the HQET, justification for fixing one of the main parameters of the model κ.

Our prediction for the branching ratio (for $|V_{cb}| = 0.040$, $\tau_{\Lambda_b} = 1.466\times10^{-12}$s) $Br^{\text{theor}}(\Lambda_b \to \Lambda_c l\nu) = (7.6 \pm 0.8)\%$ agrees within error bars with the experimental value $Br^{\text{exp}}(\Lambda_b \to \Lambda_c l\nu) = (6.2^{+1.2}_{-1.2})\%$. For other semileptonic heavy-to-heavy decays experimental data are not available at present. Comparison of different theoretical predictions can be found in Ref. [14].

In summary, we demonstrated that spectroscopy and weak decays of heavy and strange baryons can be well described in the relativistic quark model based on the quark-diquark picture of baryons.

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