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Exotic baryon resonances and the model of chiral solitons

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Contents

Abstract. A recently observed baryonic resonance with positive strangeness is discussed. The state is exotic because valence quarks alone cannot produce it and at least one quark – antiquark pair must be involved. The mass of the resonance is in agreement with chiral soliton models, which predict a number of other exotic states with different quantum numbers. Some of these states (e.g., those with strangeness $S = -2$ and charges $Q = 0, -2$) are probably already being observed experimentally.

1. Introduction

In recent experiments [1 – 4], a baryon resonance has been revealed with positive strangeness $S = +1$, charge $Q = +1$ and spin and parity that have not yet been determined. A state of mass 1.54 ± 0.01 GeV and width below 25 MeV has been observed at the confidence level 4.6σ in the missing mass spectrum of K^+n in the photoproduction reaction on carbon nuclei in experiment [1] at the SPring-8 installation.

In an ITEP experiment [2], a study of the interaction of low-energy K^+-mesons with xenon nuclei has been performed with the xenon bubble chamber DIANA, and a narrow baryon of mass 1539 ± 2 MeV and width below 9 MeV was observed in the K^0p system at the reliability level 4.4σ. The films analyzed in Ref. [2] were obtained in the 1980s. It must also be noted that the existence of a resonance in the K^0_Sp system signifies the existence of a resonance with positive strangeness only in the absence of a resonance in the $\bar{K}^0p$ system within the same mass range.

Somewhat later, in the Jefferson Laboratory, a study was performed with the large-acceptance CEBAF spectrometer (CLAS) [3] on the exclusive photoproduction reaction of K-mesons on deuterons,

$$\gamma d \to K^+K^-pn\,,$$

and a narrow resonance was revealed in the K^+n system. The mass of this state is 1542 ± 5 MeV, its width is less than 21 MeV, and the reliability level is $5.3 \pm 0.5\sigma$.

After the results of experiments [1, 2] already became known, a peak in the distribution of invariant masses in the nK^+ system corresponding to the mass $1540 \pm 4 \pm 2$ MeV was observed at a 4.8σ reliability level [4] in the $K\bar{K}$ pair production reaction with the aid of the detector SAPHIR at the Bonn electron accelerator ELSA. The upper limit for the width amounts to 25 MeV. The absence of the double-charge state Θ^{++} in the $\gamma p \to pK^+K^-$ channel has permitted the authors to draw the conclusion that the resonance Θ^+ is most likely an isoscalar.

Analysis of data obtained previously on the interaction of neutrinos and antineutrinos with nuclei in bubble chambers at CERN and at the Fermi laboratory [5] has also revealed a resonance in the K^0_Sp system with the mass 1533 ± 5 MeV and a width below 20 MeV at the reliability level 6.7σ.†

Finally, the NA49 collaboration at CERN has most recently observed resonances in the $\Xi^-\pi^-$ (charge -2) and $\Xi^-\pi^+$(charge zero) systems of mass 1.86 GeV and width below the experimental resolution, amounting to 18 MeV, at the reliability level 4.0σ [6] in proton – proton collisions at the energy 17.2 GeV in the center-of-mass system. These resonances are, most likely, components of the isotopic quartet of $\Xi^*_{3/2}$ hyperons. It is of particular interest that the same experiment revealed antibaryon resonances in the $\bar{\Xi}^+\pi^+$ and $\bar{\Xi}^+\pi^-$ systems that are the respective antiparticles of $\Xi^*_{3/2}$. A joint analysis of the data on baryons and antibaryons has led to the mass values $M_{\Xi^*_{3/2}}(Q = -2) = 1.862 \pm 0.002$ GeV and $M_{\Xi^*_{3/2}}(Q = 0) = 1.864 \pm 0.005$ GeV.

If the results of Refs [1 – 5] are confirmed (which practically no one doubts, because the resonance has been observed at different installations in Japan, Russia, the USA, and Germany, as well as at CERN), this will mean that for the

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† As was noted by A Dolgolenko, the ambiguity in the sign of strangeness of the resonance in the K^0_Sp system existing in Ref. [5] does not exist in Ref. [2] because of the low energy of K^+ mesons interacting with Xe nuclei. (*Author's note to English translation.*)

first time a baryon resonance of positive strangeness has been observed, which together with the $\Lambda(1520)$ and $\Xi(1530)$ hyperons exhibits one of the most narrow widths among all the known baryon resonances. Nor is a width of the new resonance below 9 MeV [2] excluded by the results of Refs [1, 3–5], and in such a case it has a record small width for baryon resonances. This baryon, predicted earlier theoretically in Refs [7–9], called Z^+ [9] at the beginning and then Θ^+, necessarily contains a strange antiquark $\bar{s}$, because baryons consisting of only three valence quarks can only have negative strangeness, $S < 0$ (or, naturally, zero strangeness, like the nucleon and the isobar). Thus, in addition to three valence quarks, the Θ^+ baryon contains at least one quark–antiquark pair and is therefore exotic. The resonance with strangeness -2 revealed in Ref. [6] must also contain at least one quark–antiquark pair, because the electric charge of a baryon with strangeness $S = -2$ cannot be -2 if it consists of only three valence quarks (see also the next section). Therefore, the baryon resonance found in Ref. [6] represents the second example of an experimentally established exotic baryon state. A general review of the experimental situation and of the methods applied for observing exotic hadrons as well as so-called cryptoexotic states, i.e., states with hidden exotics, is to be found, for instance, in Ref. [10].

With the same CLAS installation [3], an indication has been obtained in the $\pi^+\pi^-$ pair electroproduction reaction on protons of the existence of a new resonance with zero strangeness and positive parity, which has strong coupling with the $\Delta\pi$ channel and weak coupling with $N\rho$ and has a mass close to 1720 MeV [11]. This resonance may also belong to one of the multiplets of the exotic baryons examined in Ref. [12].

The purpose of the present article, in addition to presenting information on the revelation of essentially new resonances, is to give a short review of the theoretical ideas that led to the prediction of resonances of a new type and of the results that follow from the chiral soliton approach. As a rule, no technical details are presented, because many of the results obtained are quite transparent (we leave these details for a detailed article intended for theoreticians).

In Section 2, the main physical ideas underlying the model are explained and it is shown how multiplets of exotic baryons arise within the framework of the chiral soliton approach.

The mass formula for baryon states is presented in Section 3, its origin is explained, and we note that this formula is quite natural. In Section 4, the spectra of baryon multiplets are presented, taking into account the mixing of configurations for several versions of the model that differ in the numerical values of the parameters. In the concluding sections, some examples are discussed, a number of states whose search is of primary interest are indicated, and the results are compared with those obtained within the framework of other models.

2. Multiplets of exotic baryons in topological soliton models

Exotic baryon states, in the special meaning of this term, include states that cannot be composed of $3B$ valence quarks (B represents the baryon number) and that, by virtue of their quantum numbers, must contain one (or more) quark–antiquark pairs. Evidently, any state with positive strangeness is exotic, as is any state with a sufficiently large negative strangeness $S < -3B$. Moreover, at any value of hypercharge or strangeness $S < 0$, a state with a sufficiently large isotopic spin $I > (3B + S)/2$ is exotic. This is due only to nonstrange quarks having nonzero isotopic spin, while the number of valence nonstrange quarks is limited to the value $3B + S$. The observed hyperon of strangeness $S = +1$, containing at least one quark–antiquark pair, is also called a pentaquark state. It is well known that any baryons (generally speaking, hadrons), which, given their quantum numbers, can only consist of valence quarks, contain so-called 'sea' quarks, as well as gluons carrying away a significant part of the baryon momentum. In pentaquarks of the Θ^+ type, however, the $q\bar{q}$ pair is a valence pair, meaning that it carries certain quantum numbers, for instance, strangeness.

From the standpoint of theory, the existence of such states is not unexpected. The necessity of the existence of a baryon antidecuplet was noted by a number of authors within the framework of both quark models [13] and chiral soliton models [14, 15]. The first numerical estimations of the masses of baryon decuplet components were made in Refs [7, 15, 8]. In Refs [7, 15, 8, 9, 16], relatively small masses were predicted for states of the baryon antidecuplet, and in Refs [7, 15], there were, strictly speaking, insufficient grounds for doing so, because the mass splitting within the baryon octet and decuplet was not described. In Ref. [9], the assumption that the nucleonic resonance $N^*(1{,}71)$ represents a component of the antidecuplet is essential for estimating the mass of the state with positive strangeness, $M_\Theta = 1530$ MeV, in agreement with the estimation made earlier [7]. The width Γ_Θ of the Θ^+ hyperon was estimated in Refs [9, 16], and the result obtained in Ref. [16] is several times larger than the result in Ref. [9].

Topological soliton models are very economical and efficient in predicting the spectra of baryons or of baryon systems with various quantum numbers. Relativistic multiparticle problems, such as finding bound states and resonances in systems of three, five, etc. quarks, can certainly not be resolved within the framework of this approach; however, it turns out to be possible to calculate the spectra of baryon states without going into detail about their internal structure. In these models, baryons or baryon systems (nuclei) originate as quantized classical field configurations corresponding to the minimum of classical energy (mass). In a relatively short article, it is difficult to fully justify and explain the model of chiral solitons, and we only present essentially important explanations. The starting point is the Lagrangian of the model, written in terms of chiral fields (three pion fields for the SU(2) version of the model) and chiral derivatives, left or right, $l_\mu = \partial_\mu U U^\dagger$, $r_\mu = U^\dagger \partial_\mu U$, and $U \in \mathrm{SU}(2)$ is a unitary matrix depending on the chiral fields; in the generalization of the model in which we are interested, $U \in \mathrm{SU}(3)$. The effective Lagrangian (its space density) that describes meson interactions at low energies essentially represents an expansion in powers of chiral derivatives,

$$\tilde{L} = \tilde{L}^{(2)} + \tilde{L}_A^{(4)} + \ldots, \tag{1}$$

where

$$\tilde{L}^{(2)} = -\frac{F_\pi^2}{16}\,\mathrm{Tr}\,(l_\mu l_\mu) \tag{2}$$

is a model-independent second-order term, depending on the pion decay constant F_π, and

$$\tilde{L}_A^{(4)} = \frac{1}{32e^2}\,\mathrm{Tr}\,(l_{\mu\nu})^2 \tag{3}$$

is a soliton stabilizing the fourth-order term in the derivatives, or a Skyrme term depending on the parameter e, the Skyrme constant; $l_{\mu\nu} = \partial_\mu l_\nu - \partial_\nu l_\mu = l_\nu l_\mu - l_\mu l_\nu$ may be called the chiral field strength. There are also higher-order terms in the chiral derivatives; however, we here only deal with the version of the model with the indicated second- and fourth-order terms in the Lagrangian. It is essential that because $F_\pi^2 \sim 1/e^2 \sim N_c$, the effective Lagrangian is proportional to N_c, the number of colors of quantum chromodynamics, on the basis of which the effective theory is constructed.

In addition to chirally invariant terms in the Lagrangian depending only on chiral derivatives, there are also terms violating chiral invariance and depending on the pion field and its mass,

$$\tilde{L}_{\mathrm{CSB}} = \frac{F_\pi^2 m_\pi^2}{16}\, \mathrm{Tr}\,(U + U^\dagger - 2)\,. \tag{4}$$

This contribution is usually taken into account in calculating the functional of static energy and in its minimization. The contributions to the Lagrangian that violate symmetry with respect to flavors depend on the K-meson mass and its decay constant F_{K},

$$\tilde{L}_{\mathrm{FSB}} = \frac{F_{\mathrm{K}}^2 m_{\mathrm{K}}^2 - F_\pi^2 m_\pi^2}{24}\, \mathrm{Tr}\,(1 - \sqrt{3}\,\lambda_8)(U + U^\dagger - 2) - \frac{F_{\mathrm{K}}^2 - F_\pi^2}{48}\, \mathrm{Tr}\,(1 - \sqrt{3}\,\lambda_8)(U l_\mu l_\mu + l_\mu l_\mu U^\dagger)\,. \tag{5}$$

These contributions are very important in quantization of solitons in the SU(3) configuration space (see the following sections).

The field configurations are characterized by the topological number, which is also termed the mapping degree (index). For this number to exist, the field must have a sufficient number of components, at least three. The central assumption underlying the model and put forward by Skyrme back in 1961 [17] consists in that the topological characteristic of the field configuration, or the mapping index, can be identified with the baryon number B. To each point of space, there corresponds a chiral field calculated by a certain formula, such that as the argument goes around the ordinary space, its image in the isotopic [or SU(3)] space winds around the space an integer number of times, which is precisely the baryon number of the system. It is essential that the number of dimensions of the ordinary space, equal to three, coincides with the number of generators of the SU(2) group, which allows mapping the ordinary space onto the isotopic spin space, $R^3 \to \mathrm{SU}(2)$, leading to the formation of chiral solitons. Therefore, the fact that isotopic spin is such an important characteristic of baryons and, generally, of hadrons can be said to have a natural explanation in such models.

In 1983, the baryon number, postulated by Skyrme, was shown by Witten to represent the fourth component of the Noether current generated by the Wess–Zumino action [18]. This action is conveniently written as the 5-dimensional differential form

$$S_{\mathrm{WZ}} = \frac{-\mathrm{i}N_c}{240\pi^2} \int_\Omega \mathrm{d}^5x\, \epsilon^{\mu\nu\rho\sigma\tau}\, \mathrm{Tr}\,(l_\mu l_\nu l_\rho l_\sigma l_\tau)\,, \tag{6}$$

where Ω is a 5-dimensional domain whose boundary is 4-dimensional space-time and l_μ is the 5-dimensional generalization of l_μ in the 4-dimensional case. The density of the vector Noether current corresponding to the transformation $U \to U \exp(\mathrm{i}\phi/N_c)$ is

$$B_\alpha = \frac{1}{24\pi^2}\, \epsilon_{\alpha\beta\gamma\mu}\, \mathrm{Tr}\,(l_\beta l_\gamma l_\mu)\,. \tag{7}$$

In the general case of SU(2) skyrmions, the unitary matrix U is parameterized as $U = c_f + \mathrm{i}s_f \boldsymbol{\tau}\mathbf{n}$, the components of the unit vector $\mathbf{n}$ are $n_3 = c_\alpha$, $n_1 = s_\alpha c_\beta$, $n_2 = s_\alpha s_\beta$, $c_f = \cos f$, $s_f = \sin f$, $c_\alpha = \cos\alpha$, etc., i.e., the field configuration depends on three functions of three coordinates, f, α, and β. The baryon (topological) number is then given by

$$B = \frac{-1}{2\pi^2} \int s_f^2 s_\alpha\, I\left[\frac{(f,\alpha,\beta)}{(x,y,z)}\right] \mathrm{d}^3r\,, \tag{8}$$

where $I[(f,\alpha,\beta)/(x,y,z)] = \epsilon^{ikl}\,\partial_i f\,\partial_k\alpha\,\partial_l\beta$ is the Jacobian of the transformation from the variables (x,y,z) to the variables (f,α,β). Because the surface element of the unit 3-dimensional sphere is $\mathrm{d}S^3 = s_f^2\,\mathrm{d}f\, s_\alpha\,\mathrm{d}\alpha\,\mathrm{d}\beta$ and the entire surface equals $2\pi^2$, expression (8) shows how many times the three-dimensional sphere is winded around in integrating over 3-dimensional space, which is precisely the $R^3 \to S^3$ mapping index [SU(2) is homeomorphic to the three-dimensional sphere S^3].

The next step consists in searching for static configurations for each value of the topological (baryon) number. It has been established that the classical energy achieves a minimum in the case of configurations of the 'hedgehog' type, for which the chiral field at each point of space can be directed along the radius vector connecting the center of the soliton and the point of observation. For such configurations, $\mathbf{n} = \mathbf{r}/r$ and the profile function $f = f(r)$ depends only on a single variable — the distance from the center of the skyrmion. The baryon number is then given by

$$B = \frac{1}{\pi}\left[f(0) - f(\infty)\right], \tag{9}$$

and is therefore determined by the difference between the values of the profile function at the center of the skyrmion and at infinity. For baryons, one usually assumes $f(0) = \pi$ and $f(\infty) = 0$. Then the static energy functional (mass) of the skyrmion is also significantly simplified, and therefore the profile function can be calculated numerically with a computer in several seconds. For other baryon numbers, the field configurations corresponding to an energy minimum exhibit different shapes, and the distributions of mass and of the baryon number density exhibit shapes that are not spherically symmetric, for example, toroidal for $B = 2$. Simultaneously with the mass, other characteristics are also calculated, including skyrmion moments of inertia, which play a very important role in the quantization procedure. The skyrmion mass and moments of inertia are proportional to N_c, and are therefore large in the limit of a large number of colors.

In SU(3) generalizations of the model, SU(2) skyrmions are embedded in the SU(3) configuration space [in the left upper corner of the SU(3) matrix] and time-dependent matrices of collective coordinates, $A(t) \in \mathrm{SU}(3)$, are introduced such that $U_{\mathrm{SU}(3)}(\mathbf{r},t) = A(t)\,U_{\mathrm{SU}(2)}(\mathbf{r})\,A^\dagger(t)$. In the SU(3) space, the skyrmion angular velocities of rotation, ω_k, are determined by the relations

$$A^\dagger(t)\dot{A}(t) = \frac{-\mathrm{i}\omega_k\lambda_k}{2}\,, \tag{10}$$

where $k = 1, \ldots, 8$, λ_k are the Gell-Mann matrices, and the matrix $A(t)$ is usually written as

$$A = A_{\mathrm{SU}(2)} \exp(\mathrm{i}\nu\lambda_4)\, A'_{\mathrm{SU}(2)} \exp\left(\frac{\mathrm{i}\rho\lambda_8}{\sqrt{3}}\right).$$

For the $U_{\mathrm{SU}(3)}$ parameterization indicated, the Wess–Zumino–Witten action is readily calculated; moreover, it is possible to calculate the corresponding contribution to the Lagrangian [19],

$$L_{\mathrm{WZ}} = -\omega_8 \frac{N_{\mathrm{c}} B}{2\sqrt{3}}, \tag{11}$$

which depends only on ω_8.[1] The hypercharge Y_{R} of a baryon (or of a baryon system) in the reference frame attached to the rotating soliton, also termed the 'right' hypercharge, is determined by

$$Y_{\mathrm{R}} = -\frac{2}{\sqrt{3}} \frac{\partial L}{\partial \omega_8}.$$

Hence arises the so-called Guadagnini quantization condition [19]

$$Y_{\mathrm{R}} = \frac{N_{\mathrm{c}} B}{3}. \tag{12}$$

For any SU(3) multiplet (p, q), the maximum hypercharge, or triality, $Y_{\max} = (p + 2q)/3$, and the inequality

$$\frac{p + 2q}{3} \geqslant \frac{N_{\mathrm{c}} B}{3} \tag{13}$$

must evidently be satisfied; therefore,

$$p + 2q = 3(B + m) \tag{14}$$

for $N_{\mathrm{c}} = 3$, where m is a positive integer. This quantization condition has a simple physical interpretation: we proceed from an initially nonstrange configuration, which remains such in the system attached to the soliton. All other states of the (p, q) multiplet arise in the reference system of the observer as a result of this configuration rotating in the SU(3) configuration space and are correspondingly described by the Wigner functions of an SU(3) rotating top, and each multiplet must contain a nonstrange state. States with $m = 0$ are, naturally, called minimal states [21]; for $B = 1$, the octet and decuplet are minimal multiplets, while multiplets of lower dimensions are forbidden by the Guadagnini condition [19] [we recall that the number of components of a multiplet is $N(p, q) = (p + 1)(q + 1)(p + q + 2)/2$].

States with $m = 1$ contain at least one quark–antiquark pair. Indeed, the maximum hypercharge is 2, which corresponds to strangeness $S = +1$, i.e., a $\mathrm{q}\bar{\mathrm{s}}$ pair should be present, q = u or d, and by virtue of the SU(3) invariance of strong interactions, all components of the multiplet must contain a $\mathrm{q}\bar{\mathrm{q}}$ pair. An additional restriction arises when isotopic spin is considered.

[1] For the parameterization $U_{\mathrm{SU}(3)} = A(t)U'(\mathbf{r})A^{\dagger}(t)$, the contribution of the Wess–Zumino–Witten action to the Lagrangian is always linear in the angular velocities ω_k; however, for arbitrary $U' \in \mathrm{SU}(3)$, all the eight components of the angular velocities may contribute. The general case of the SU(3)-skyrmion quantization is considered in Ref. [20].

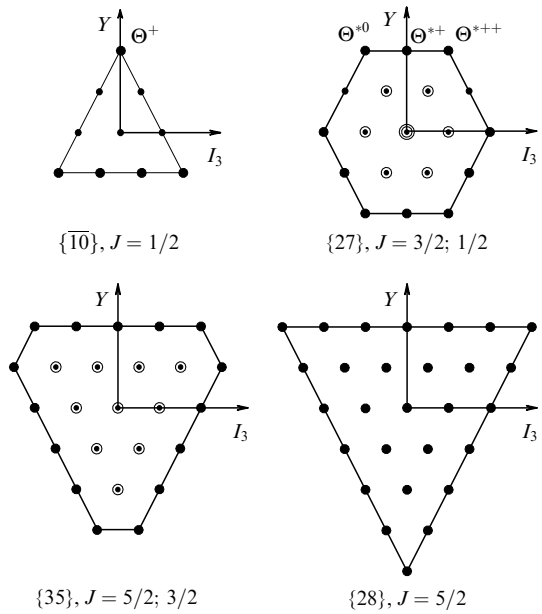

Figure 1. $I_3 - Y$ diagrams of $B = 1$, $m = 1$ baryon multiplets. The large black circles indicate exotic states, the small ones indicate cryptoexotic states, which can mix with the respective components of the baryon octet and decuplet.

It is not difficult to verify that the $\{\overline{10}\}$-, $\{27\}$-, and $\{35\}$-plets are pentaquark states, but the state with maximum p, $(p, q) = (6, 0)$, or the $\{28\}$-plet, must already contain two $\mathrm{q}\bar{\mathrm{q}}$ pairs (i.e., be a septuquark), because this multiplet contains a state of strangeness $S = -5$ and a state of strangeness $S = +1$ and isospin $I = 3$ [obviously, the pentaquark $(\mathrm{qqqq}\bar{\mathrm{s}})$ cannot have isospin greater than 2].

The $I_3 - Y$ diagrams for multiplets with $m = 1$ are depicted in Fig. 1. The minimum hypercharge in the multiplet is $Y_{\min} = -(2p + q)/3$, and the maximum isospin $I_{\max} = (p + q)/2$ is achieved at $Y = (p - q)/3$. Complex multiplets, such as $\{27\}$, $\{35\}$, for $m = 1$ and all others, with the exception of the last one with $(p, q) = (9, 0)$, for $m = 2$, contain two or more states with different values of spin J at the inner points.

3. The mass formula

In addition to the contribution to the Lagrangian linear in the angular velocities of rotation in the SU(3) space, due to the Wess–Zumino action, there exist contributions quadratic in these angular velocities that arise from the Lagrangian in Eqns (1)–(3), i.e., the rotation energy $E_{\mathrm{rot}} = L_{\mathrm{rot}}$. The coefficients in the respective quadratic form are moments of inertia, isotopic or pion Θ_π and strange (flavor) or kaon Θ_{K} [19]:

$$L_{\mathrm{rot}} = \frac{1}{2}\,\Theta_\pi(\omega_1^2 + \omega_2^2 + \omega_3^2) + \frac{1}{2}\,\Theta_{\mathrm{K}}(\omega_4^2 + \ldots + \omega_7^2). \tag{15}$$

Below, we give not very unwieldy formulas for these moments of inertia, expressed through the skyrmion profile functions.

Expression (15) is obtained quite rigorously from the initial Lagrangian of the model, but we do not go into technical details here, because the physical results obtained are clear. It is essential that for the embedding of SU(2) into SU(3) under consideration, the rotation energy is indepen-

dent of ω_8 (λ_8 and the initial ansatz commute with each other).

The Hamiltonian arises from (15) when the canonical quantization procedure [19, 8, 9, 12] is applied,

$$H = M_{\rm cl} + \frac{1}{2\Theta_\pi}\,\mathbf{R}^2 + \frac{1}{2\Theta_{\rm K}}\left[C_2({\rm SU}_3) - \mathbf{R}^2 - \frac{N_{\rm c}^2 B^2}{12}\right], \quad (16)$$

where the second-order Casimir operator for the SU(3) group is

$$C_2({\rm SU}_3) = \sum_{a=1}^{8} R_a^2$$

and the eigenvalues for the (p,q)-multiplets are

$$C_2({\rm SU}_3)_{p,q} = \frac{p^2+pq+q^2}{3} + p + q\,,$$

and

$$\mathbf{R}^2 = R_1^2 + R_2^2 + R_3^2 = J(J+1) = I_{\rm R}(I_{\rm R}+1)$$

for the SU(2) group.

The operators $R_a = \partial L/\partial\omega_a$ satisfy certain commutation relations that represent an SU(3) generalization of the commutation relations for components of the angular momentum [19]. The terms in the Lagrangian that are linear in the rotation angular velocity drop out from Hamiltonian (16). For minimal multiplets ($m=0$), the so-called right isospin is $I_{\rm R} = p/2$. It is readily verified that for such multiplets, the coefficient at $1/2\Theta_{\rm K}$ is

$$K = C_2({\rm SU}_3) - \mathbf{R}^2 - \frac{3B^2}{4} = \frac{3B}{2}\,, \quad (17)$$

and is therefore the same for all minimal multiplets ($m=0$), $N_{\rm c}=3$ [21] (Table 1).[2]

Table 1. Values of $N(p,q)$, of the Casimir operator $C_2({\rm SU}_3)$, of the right isospin $I_{\rm R}$, and of the coefficients $K(J_{\max})$ and $K(J_{\max}-1)$ for the minimal and first nonminimal baryon multiplets.

(p,q)	$N(p,q)$	m	$C_2({\rm SU}_3)$	$J = I_{\rm R}$	$K(J_{\max})$	$K(J_{\max}-1)$
(1, 1)	{8}	0	3	1/2	3/2	
(3, 0)	{10}	0	6	3/2	3/2	
(0, 3)	$\{\overline{10}\}$	1	6	1/2	3/2 + 3	
(2, 2)	{27}	1	8	3/2; 1/2	3/2 + 2	3/2 + 5
(4, 1)	{35}	1	12	5/2; 3/2	3/2 + 1	3/2 + 6
(6, 0)	{28}	1	18	5/2	3/2 + 7	
(1, 4)	$\{\overline{35}\}$	2	12	3/2; 1/2	3/2 + 6	3/2 + 9
(3, 3)	{64}	2	15	5/2; 3/2; 1/2	3/2 + 4	3/2 + 9
(5, 2)	{81}	2	20	7/2; 5/2; 3/2	3/2 + 2	3/2 + 9
(7, 1)	{80}	2	27	7/2; 5/2	3/2 + 9	3/2 + 16
(9, 0)	{55}	2	36	7/2	3/2 + 18	

For $B=1$, this property of the baryon mass spectrum was already noted in Ref. [19] [see (18)]. For nonminimal multiplets, there arise additional contributions to the energy, $\delta E \sim m/\Theta_{\rm K}$ and $\sim m^2/\Theta_{\rm K}$ [21] [see (19)–(22)]. Thus, within the framework of the chiral soliton approach, the 'weight' of a quark–antiquark pair is determined by the parameter $1/\Theta_{\rm K}$, and this remarkable property of the model merits a more profound understanding.

The equality of the numerical values of the angular momentum J and of the isotopic spin $I_{\rm R}$ in the reference system attached to the soliton (of the body-fixed, or 'right' isospin), applied in (15) and (16), is valid only for configurations of the 'hedgehog' type, for which space and isotopic rotations are equivalent, and is not valid for other configurations realizing an energy minimum for $B \geqslant 2$.

From Table 1, it follows that at each m, the coefficient $K(J_{\max})$ decreases as $N(p,q)$ increases, for example, $K(35) < K(27) < K(\overline{10})$. It is not difficult to obtain the following values for the contributions of rotational energy (16) to the mass differences of multiplets:

$$M_{\{10\}} - M_{\{8\}} = \frac{3}{2\Theta_\pi} \quad (18)$$

(this relation has been known since 1984 [19]),

$$M_{\{\overline{10}\}} - M_{\{8\}} = \frac{3}{2\Theta_{\rm K}}\,, \quad (19)$$

which was especially stressed in Ref. [9],

$$M_{\{27\},J=3/2} - M_{\{10\}} = \frac{1}{\Theta_{\rm K}}\,, \quad (20)$$

$$M_{\{27\},J=3/2} - M_{\{\overline{10}\}} = \frac{3}{2\Theta_\pi} - \frac{1}{2\Theta_{\rm K}}\,, \quad (21)$$

$$M_{\{35\},J=5/2} - M_{\{27\},J=3/2} = \frac{5}{2\Theta_\pi} - \frac{1}{2\Theta_{\rm K}}\,. \quad (22)$$

Therefore, if the relation $\Theta_{\rm K} \ll \Theta_\pi$ were valid, the {27}-plet would be lighter than the antidecuplet and the {35}-plet would be lighter than the {27}-plet. In reality, $\Theta_{\rm K}$ is approximately two times smaller than Θ_π, and hence the components of the antidecuplet turn out to be lighter than the corresponding (in strangeness) components of the {27}-plet. Starting from a certain value of $N(p,q)$, the coefficient K increases drastically, which corresponds to an increase in the number of quark–antiquark pairs by one unit (see Table 1). Because the moments of inertia are such that $\Theta_{\rm K} \sim \Theta_\pi \sim N_{\rm c}$, the energy differences of multiplets with $m=0$ $\sim 1/N_{\rm c}$ [see (18)], while the differences in the rotation energy of states with different m contain contributions $\sim N_{\rm c}/\Theta_{\rm K} \sim 1$ [see (19), (20)].

Formula (16) was obtained in the rigid-rotator approximation, i.e., with the soliton profile function and, correspondingly, its dimensions and other properties assumed unchanged under rotation in configuration space. For this approximation to be valid, it is necessary that the soliton rotation time $\tau_{\rm rot}$ in configuration space be smaller than its deformation time $\tau_{\rm deform}$ under the influence of forces due to the presence of terms violating flavor symmetry in the Lagrangian, i.e., $m_{\rm K}/m_\pi > 1$, $F_{\rm K}/F_\pi > 1$, where $F_{\rm K}$, F_π are the decay constants of K- and π-mesons, and $F_{\rm K}/F_\pi \simeq 1.22$ experimentally. The time $\tau_{\rm rot}$ can be easily estimated, because $\tau_{\rm rot} \sim \pi/\omega$ and $\omega \sim \sqrt{C_2({\rm SU}_3)}/\Theta_{\rm K}$. It is more difficult to estimate $\tau_{\rm deform}$; one can only assert that it is larger than the propagation time of the signal inside the soliton, $\tau_{\rm sign} \sim 2R_{\rm H}$.

[2] For an arbitrary number of colors, $K = N_{\rm c}B/2$, but the minimal baryon multiplets then differ from the octet and decuplet for $N_{\rm c}=3$. For example, for $N_{\rm c}=5$ and $B=1$, the following are minimal multiplets: the {15}-plet with $(p,q)=(1,2)$, the {24}-plet $[(p,q)=(3,1)]$, and the {21}-plet $[(p,q)=(5,0)]$.

Thus, the rigid-rotator approximation is valid if

$$\pi\Theta_{\mathrm{K}} \ll 2R_{\mathrm{H}}\sqrt{C_2(\mathrm{SU}_3)}\,.$$

Numerically, we obtain $\pi\Theta_{\mathrm{K}} \simeq 8$ GeV^{-1} and $2R_{\mathrm{H}}\sqrt{C_2(\mathrm{SU}_3)} \simeq 12$ GeV^{-1} for the baryon decuplet or antidecuplet.

An alternative to the above is the 'soft', or slow, rotator approximation, when it is assumed that at each rotation angle v in the 'strange' direction, there is sufficient time for the soliton to undergo deformation under the action of forces due to the existence of terms violating flavor symmetry in the Lagrangian [22]. The reality is intermediate; however, for baryons, the rigid-rotator approximation is more justified owing to the estimation made above. The soft or slow rotator becomes preferable as the baryon number increases. The dependence of the moments of inertia on v, the rotation angle of the soliton in the strange direction, is determined by the expressions [22, 23]

$$\Theta_{\mathrm{K}}(v) = \frac{1}{8}\int(1-c_f)\left[F_{\mathrm{K}}^2\left(1-\frac{2-c_f}{2}\,s_v^2\right)+F_\pi^2\,\frac{2-c_f}{2}\,s_v^2 + \frac{1}{e^2}\left(f'^2+\frac{2s_f^2}{r^2}\right)\right]\mathrm{d}^3r\,, \tag{23}$$

$$\Theta_\pi(v) = \frac{1}{6}\int s_f^2\left[F_\pi^2 + c_f(F_{\mathrm{K}}^2-F_\pi^2)s_v^2 + \frac{4}{e^2}\left(f'^2+\frac{s_f^2}{r^2}\right)\right]\mathrm{d}^3r\,. \tag{24}$$

These expressions are valid for configurations of the 'hedgehog' type, when the skyrmion is described by a sole profile function f, $s_v = \sin v$. $\Theta_{\mathrm{K}}(v)$ decreases as v increases; $\Theta_\pi(v)$, on the contrary, increases, and $v = 0$ corresponds to the rigid-rotator approximation. Actually, the truth lies somewhere in between, and for $B = 1$, the value $v = 0$ is close to the correct value.

4. The spectrum of baryon states

Expressions (15) and (16), together with the figures given in Table 1, are sufficient for calculating baryon spectra without taking the mass splitting within the multiplets into account, as in Refs [15, 21]. The mass splitting, due to the presence of terms violating flavor symmetry in the Lagrangian [see expression (5)] plays a very essential role [19, 7, 8, 12],

$$H_{\mathrm{SB}} = \frac{1-D_{88}^{(8)}}{2}\,\Gamma_{\mathrm{SB}}\,, \tag{25}$$

where the SU(3) Wigner function is $D_{88}^{(8)}(v) = 1 - 3s_v^2/2$ and

$$\Gamma_{\mathrm{SB}} = \frac{2}{3}\left[\left(\frac{F_{\mathrm{K}}^2}{F_\pi^2}\,m_{\mathrm{K}}^2 - m_\pi^2\right)\Sigma + (F_{\mathrm{K}}^2-F_\pi^2)\tilde{\Sigma}\right], \tag{26}$$

$$\Sigma = \frac{F_\pi^2}{2}\int(1-c_f)\,\mathrm{d}^3r\,, \tag{27}$$

$$\tilde{\Sigma} = \frac{1}{4}\int c_f\left(f'^2+\frac{2s_f^2}{r^2}\right)\mathrm{d}^3r\,. \tag{28}$$

The quantity $\mathrm{SC} = \langle s_v^2\rangle/2 = \langle 1 - D_{88}^{(8)}\rangle/3$ averaged over the baryon wave function represents its strangeness content. Without taking the mixing of configurations into account, i.e., when the symmetry-violating contributions to the Lagrangian are considered a small perturbation, $\langle s_v^2\rangle_0$ is simply expressed via the Clebsh–Gordan coefficients of the SU(3) group. The values of $\langle s_v^2\rangle_0$ for the octet, decuplet, antidecuplet, and certain components of higher multiplets are presented in Table 2. In this approximation, the components $\{10\}$ and $\{\overline{10}\}$ are situated equidistantly [19, 7–9], and the mass splittings of the decuplet and antidecuplet are equal to each other.

The spectrum of states can be obtained with configuration mixing taken into account by diagonalizing the Hamiltonian in the next order of the perturbation theory in H_{SB} (the

Table 2. Mass values for the octet, decuplet, antidecuplet, and certain components of higher baryon multiplets for different values of the Skyrme constant: version A — $e = 3.96$, version B — $e = 4.12$, version C — fit with parameters Θ_{K}, Γ_{SB} [12]. The nucleon mass is assumed to be equal to the observed value.

		A	B	C	
Θ_π, GeV^{-1}	—	6.175	5.556	5.61	—
Θ_{K}, GeV^{-1}	—	2.924	2.641	2.84	—
Γ_{SB}, GeV^{-1}	—	1.369	1.244	1.45	—
Baryon $\lvert N, Y, I, J\rangle$	$\langle s_v^2\rangle_0$	A	B	C	Experiment
$\Lambda\lvert 8,0,0,1/2\rangle$	0.60	1094	1078	1103	1116
$\Sigma\lvert 8,0,1,1/2\rangle$	0.73	1202	1182	1216	1193
$\Xi\lvert 8,-1,1/2,1/2\rangle$	0.80	1310	1274	1332	1318
$\Delta\lvert 10,1,3/2,3/2\rangle$	0.58	1228	1258	1253	1232
$\Sigma^*\lvert 10,0,1,3/2\rangle$	0.67	1357	1372	1391	1385
$\Xi^*\lvert 10,-1,1/2,3/2\rangle$	0.75	1483	1484	1525	1530
$\Omega\lvert 10,-2,0,3/2\rangle$	0.83	1604	1587	1654	1672
$\Theta^+\lvert\overline{10},2,0,1/2\rangle$	0.50	1519	1564	1539	1540
$N^*\lvert\overline{10},1,1/2,1/2\rangle$	0.58	1633	1664	1661	1710?
$\Sigma^*\lvert\overline{10},0,1,1/2\rangle$	0.67	1731	1749	1764	1770?
$\Xi^*_{3/2}\lvert\overline{10},-1,3/2,1/2\rangle$	0.75	1753	1781	1786	1862?
$\Theta^*\lvert 27,2,1,3/2\rangle$	0.57	1646	1697	1690	—
$\Omega^*\lvert 27,-2,1,3/2\rangle$	0.82	1930	1950	1987	—
$X\lvert 35,1,5/2,5/2\rangle$	0.44	1723	1817	1792	—
$\lvert 35,-3,1/2,5/2\rangle$	0.85	2208	2251	2306	—
$\lvert 28,2,3,5/2\rangle$	0.61	2877	3075	2982	—
$\lvert 28,-4,0,5/2\rangle$	0.78	3160	3318	3284	—

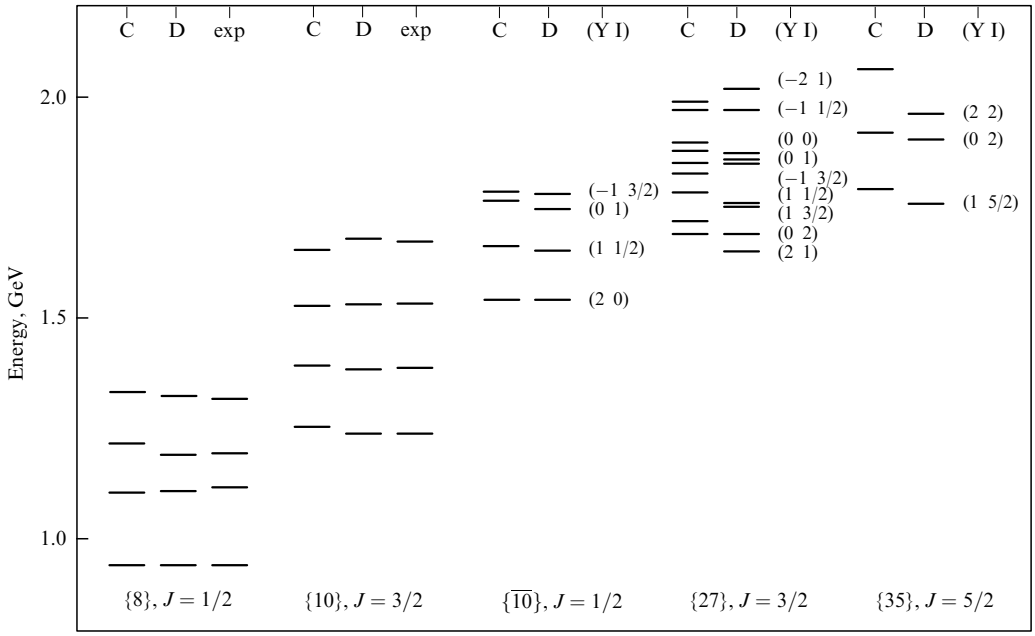

Figure 2. Baryon spectrum of the octet, decuplet, antidecuplet, {27}-plet, and certain components of the {35}-plet with configuration mixing taken into account. The figure is from Ref. [12].

computer program for calculations was made available by H Walliser). The results of calculations within the Skyrme model framework with a single parameter, the Skyrme constant e (the experimental value $F_\pi = 186$ MeV), are presented in Table 2. The equidistant position of the decuplet and, in particular, of the antidecuplet components is quite evidently violated when configuration mixing is taken into account. As a rule, the quantity $\langle s_\nu^2 \rangle$ decreases when configuration mixing is taken into account.

It is necessary to bear in mind that in its modern state, the chiral soliton approach only describes the mass differences of baryons or of baryon systems [8, 9, 12, 22]. The absolute mass values are controlled by loop corrections $\sim N_c^0 \sim 1$, which at present have been roughly estimated only for solitons with $B = 1$ [24]. Therefore, the values of the nucleon mass given in Table 2 and Fig. 2 are set equal to the observed values.

From Table 2, one can see that the agreement with data on the masses of the baryon octet and decuplet for versions A and B is only satisfactory on the whole; however, the observed mass of the Θ^+ hyperon is obtained without difficulty and even with a certain reserve. For a more reliable prediction of the masses of other exotic baryons, a phenomenological approach was adopted in Ref. [12], in which the measured baryon mass $M_{\Theta^+} = 1.54$ GeV was used as the input value and a mass formula was applied with parameters Θ_K, Θ_π, and Γ_{SB}. The results thus obtained are presented in Table 2 (version C) and in Fig. 2. The masses of some components of the {27}-, {35}-, and {28}-plets are also given. Version D in Fig. 2 corresponds to the case that accounts for the term in H_{SB} that arises, for example, from $\rho-\omega$-meson mixing in the effective Lagrangian of the model [8, 12],

$$H_{SB}^{(2)} = -\frac{\Delta}{\Theta_\pi} \sum_{a=1}^{3} D_{8a} R_a \,. \tag{29}$$

The best description of the data on the baryon masses of the octet and decuplet is achieved at $\Delta = 0.4$. Such a contribution to H_{SB} was also taken into account in Ref. [9], where a very important role is attributed to the term $H_{SB}^{(Y)} = \beta Y$ in the Hamiltonian, also linear in the hypercharge, with $\beta \simeq -156$ MeV, which is absent in the approach of Refs [8, 12].

At first sight, it may seem surprising that the state Θ^+, containing a strange antiquark, turned out to be lighter than the nonstrange component of the antidecuplet $N^*(I = 1/2)$. However, this is readily explained if one considers that all the components of the antidecuplet contain a $q\bar{q}$ pair: Θ^+ contains four light quarks and an antiquark $\bar{s}$; N^* contains three light quarks and a $q\bar{q}$ pair, including an $s\bar{s}$ pair; and $\Sigma^* \in \{\overline{10}\}$ contains u, d, and s quarks and an $s\bar{s}$ pair, and so on.

The mass splitting within the baryon decuplet is significantly affected by its mixing with components of the {27}-plet, which considerably enhances the splitting [12] — an effect not taken into account in Ref. [9]. Mixing of the antidecuplet with the baryon octet leads to a certain change in the position of the components N^* and Σ^*, while the masses of Θ^+ and $\Xi^*_{3/2}$ remain unchanged: they are influenced by the mixing with multiplets of higher dimensions [12].

The component of the {35}-plet with zero strangeness deserves attention, because it has the smallest value of s_ν^2, smaller than that of the nucleon and of the Δ-isobar. Owing to the isospin conservation by strong interactions, this resonance can decay into $\Delta\pi$, but not into $N\pi$ or $N\rho$. In accordance with the results presented in Table 1, the components of the {28}-plet have masses significantly exceeding the masses of other multiplets with $m = 1$.

All the discussed baryon states, obtained from quantization of soliton rotations in the SU(2) or SU(3) configuration space, have the same, i.e., positive, parity. Qualitative discussions of the influence of other (nonzero) modes, vibrational and respiratory, and references to the literature on these issues can be found, e.g., in Refs [12, 16]. The actual

situation may be much more complicated than the simplified picture presented here because each rotational state may have several vibrational excitations with a characteristic energy up to several hundred megaelectronvolts.

If the matrix element of the $\Theta^+ \to \mathrm{KN}$ decay is represented as

$$M_{\Theta^+\to\mathrm{KN}} = g_{\Theta\mathrm{KN}}\bar{u}_{\mathrm{N}}\gamma_5 u_\Theta \phi_{\mathrm{K}}^\dagger\,, \tag{30}$$

where u_Θ and u_{N} are bispinors of the initial and final baryons, then the decay width is

$$\Gamma_{\Theta\to\mathrm{KN}} = \frac{g^2_{\Theta\mathrm{KN}}}{8\pi}\,\frac{\Delta_M^2 - m_{\mathrm{K}}^2}{M^2}\,p_{\mathrm{K}}^{\mathrm{c.m}}\,, \tag{31}$$

where M is the mass of the decaying baryon, $\Delta_M = M - m_{\mathrm{N}}$, and $p_{\mathrm{K}}^{\mathrm{c.m}}$ is the kaon momentum in the center-of-mass reference system. One can approximately, but with very good precision, write

$$\Gamma_{\Theta\to\mathrm{KN}} \simeq \frac{g^2_{\Theta\mathrm{KN}}}{8\pi}\,\frac{(p_{\mathrm{K}}^{\mathrm{c.m}})^3}{m_{\mathrm{N}}M}\,. \tag{32}$$

Here, the momentum dependence of the decay probability, characteristic of the p-wave, is singled out explicitly. For the decay constant, we obtain $g_{\Theta\mathrm{KN}} \simeq 4.4$ if the experimental value $\Gamma_{\Theta\to\mathrm{KN}} = 10$ MeV is adopted. This value is to be compared with the value $g_{\pi\mathrm{NN}} \simeq 13.5$ or with the value $g_{\pi\Xi^*\Xi} \sim 7$ for the decay $\Xi^*(1530) \to \Xi\pi$. Thus, suppression of the $\Theta^+ \to \mathrm{KN}$ decay does occur, but it is not very large and qualitatively explicable, according to the authors of Refs [25, 36]. The width of Θ^+ obtained in Ref. [9] amounts to less than 17.5 MeV, with the influence of the factor $(p_{\mathrm{K}}^{\mathrm{c.m}})^3$ on the decay probability taken into account.[3] According to a recent calculation [26], the presence of the Θ^+ resonance does not contradict the phase analysis of data on KN-scattering at energies in the 1520 – 1560 MeV interval if its width does not exceed ~ 1 MeV.

5. Discussion, prospects

The mass of the recently revealed baryon Θ^+ with positive strangeness approximately complies with predictions made earlier within the topological (chiral) soliton model [7 – 9, 16]. The width Γ_Θ calculated in Refs [9, 16] may exceed the observed value. If, on the contrary, it turns out to be that $\Gamma_\Theta \sim 1$ MeV, as required by phase analysis [26], then this may present a serious problem for the chiral soliton approach and for other theoretical models.

We also note that the antidecuplet mass splitting obtained in Refs [8, 12] is significantly smaller than in Ref. [9], where it amounted to 540 MeV. Moreover, the deviation from equidistance, resulting from configuration mixing being taken into account more completely, is significant [8, 12]. As a result, the mass of the isospin-3/2 hyperon $\Xi^*_{3/2}$ obtained in Ref. [12] is significantly smaller than in Ref. [9], and amounts to 1790 MeV, to be compared to 2070 MeV in Ref. [9]. It is interesting to note that the estimation of the $\Xi^*_{3/2}$ mass made in Ref. [25] within the quark model yielded 1750 MeV, which is close to our result [12]. The mass value of about 1860 MeV obtained experimentally [6] for the resonance in the $\Xi\pi$ system is the value closest to the prediction in Ref. [12]. The mass of $\Sigma^* \in \{\overline{10}\}$ is also lower than in Ref. [9] and is closer to the $\Sigma^*(1770)$ mass than to the $\Sigma^*(1880)$ mass.

To verify that Θ^+ does indeed belong to the antidecuplet, it is necessary, first of all, to determine its spin and parity, as well as the masses of its partners in the SU(3)-multiplet. From this point of view, the experimental results in Ref. [6] are quite promising, signifying that the partner of Θ^+ in the antidecuplet with the largest (in absolute value) strangeness is most likely observed. It remains to find the components of this quartet with charges -1 and $+1$ that are lacking.

Of interest are searches of the state $\Theta^* \in \{27\}$ with isospin $I = 1$. Correspondingly, the double-charged state Θ^{*++} may be manifested as a resonance in the $\mathrm{K}^+\mathrm{p}$ system. Because this state is 120 – 160 MeV higher than Θ^+ [12], its width should be expected to be at least 3 – 4 times larger than the Θ^+ width. The absence of such a resonance in experiments may present a serious problem for the entire chiral soliton approach.

Probably, in experiments such as those in Ref. [6], searches for resonances in the $\Xi\bar{\mathrm{K}}$ system with strangeness $S = -3$ are possible. The model of chiral solitons predicts a state with isospin $I = 1$, belonging to the $\{27\}$-plet, with charges $Q = -2, -1$, and 0 and mass close to 2000 MeV (see Table 2).

Searching for the isodoublet with strangeness $S = -4$, belonging to the $\{35\}$-plet, requires analysis of distributions in the $\Xi\bar{\mathrm{K}}\bar{\mathrm{K}}$ system or in the $\Omega\bar{\mathrm{K}}$ system, which is certainly much more difficult.

A number of exotic resonances of interest have large values of isospin, and cannot therefore be revealed in reactions of pion or kaon scattering on nucleons, but can be seen in reactions in which two or more mesons are produced, similar to the reaction studied in Ref. [11]. The possibility of identifying the state of mass ~ 1.72 GeV, found in Ref. [11], with the exotic component of the $\{35\}$-plet $S = 0$, $I = 5/2$, seems attractive (see Table 2). However, the isospin selection rule for electroproduction on a proton involving one photon exchange [11] renders this interpretation quite improbable. Another possibility, indicated by the authors of Ref. [11], is the cryptoexotic component of the $\{27\}$-plet with isospin 3/2 and mass about 1.76 GeV in accordance with the calculations in Ref. [12].

It must be noted that, naturally, there exists no contradiction between the chiral soliton approach and the quark picture of baryons and baryon resonances [25, 27], as asserted, for instance, in Ref. [27]. These approaches are complementary. The first describes baryons or baryon systems, if one can say so, from the standpoint of large distances and allows calculating those characteristics for which details of the internal baryon structure are not important. Such characteristics of baryons include, first of all, their mass. Naturally, there exist other properties of resonances for which the details of the internal structure may be more important, for instance, their electroweak properties and partial decay widths. An important consequence of the observation of pentaquark states, predicted within the framework of chiral (topological) soliton models may be a change in attitude towards such models and greater confidence in their predictions, for instance, in the case of multibaryon systems.

[3] It must be noted that in Ref. [16], a significantly larger Θ^+ decay width, close to 70 MeV if normalized to the mass $M_\Theta = 1540$ MeV, was obtained. In the same work, an error was noted in the calculations in Ref. [9]: actually, the decay width obtained in Ref. [9] exceeds the value presented by the authors. Therefore, suppression of the Θ^+ decay probability within the framework of the soliton approach cannot be considered explained, in the opinion of the authors of Ref. [16].

In spite of certain differences in the mass values and, especially, the widths of exotic resonances obtained in various versions of the model by different authors, it is an indisputable fact that chiral soliton models have allowed predicting values close to the ones subsequently measured in experiments. One can hope that the results in Refs [1 – 6, 11] open an interesting new chapter in the physics of baryon resonances and that those to follow can also be devoted to baryon systems (nuclei) with exotic properties. For example, the mass value obtained in Ref. [23] for the dibaryon with positive strangeness and isospin 1/2, $D'_{S=+1} \in \{\overline{35}\}$, is about 600 MeV higher than the NN-scattering state (singlet deuteron). Calculations performed within the framework of the bound-state model show that the excitation energy of antistrangeness (antiflavor in the general case) does not increase with the baryon number of multiskyrmions [28], and therefore searches for baryon systems with positive strangeness, as well as with heavy 'antiflavors' (anticharm, antibeauty) are important. The successful discovery of pentaquark states permits us to adopt a more serious attitude towards these predictions. The possible existence of low-lying super-narrow dibaryons of proper (electromagnetic) width ~ 1 keV below the NNπ threshold is quite interesting [28, 29]. Such dibaryons have been observed in some experiments [30, 31], but not in another experiment [32] at a mass value below 1914 MeV. Further searches are of great interest, since such a dibaryon essentially represents an elementary particle with baryon number $B = 2$.

6. Conclusion

Experiments [1 – 6, 11] provide a stimulus for the development of a new and extremely interesting sector of hadron spectroscopy — the spectroscopy of exotic baryons. The discovery of pentaquarks is certainly one of the most significant events in elementary particle physics of the past two decades. Of great interest are searches, already under way, for the missing components of the baryon antidecuplet, as well as of the components of higher multiplets ({27}-, {35}-plet, etc.). Phase analysis of data on kaon scattering on nucleons within the broadest possible range of energies and, in particular, tests of the data themselves, could yield important and interesting results.

An important, although more complicated, task consists in searching for pentaquark states containing anticharm and antibeauty, i.e., having quark contents (uudd$\bar{c}$) and (uudd$\bar{b}$). Estimates obtained within the soliton model show that such baryons may be bound with respect to decays into a nucleon and a $\bar{D}$- or a B-meson [33]. In this case, such pentaquarks will have lifetimes characteristic of the weak interaction.

Recently, a number of publications have appeared in which exotic baryon resonances in the quark model, in chiral soliton models, and within the framework of other approaches are discussed [34 – 41]; a critical review of many of them can be found in Ref. [39]. Of significant interest is report [41], in which, for instance, it is shown that in the quark model, the total mass splitting in the antidecuplet is about three times smaller than the mass splitting in the baryon decuplet. Some papers are devoted to calculations of the cross sections of photo-, electro-, and hadro-production of exotic baryons; however, this topic is not considered in the present article. Information on the Workshop held in November 2003 at the Jefferson Laboratory (CEBAF) devoted to pentaquark states can be found at the address www.jlab.org/intralab/calendar/archive03/pentaquark.

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Comments in proof

(1) The existence of an exotic baryon resonance with positive strangeness has also been confirmed by the collaboration HERMES, DESY (hep-ex/0312044) in the photoproduction reaction of the $K^0_S p$ system on deuterons. The mass of this state with $|S| = 1$ amounts to $1528 \pm 2.6 \pm 2.1$ MeV at a $4-6\sigma$ confidence level, while its width exceeds 4.3 – 6.2 MeV. The CLAS collaboration has confirmed the existence of Θ^+ in the photoproduction reaction on protons, $\gamma p \to \pi^+ K^- K^+ n$ (hep-ex/0311046). The mass of the resonance in the K^+n system amounted to 1555 ± 10 MeV at a $7.8 \pm 1\sigma$ confidence level. Critical analysis of the experimental situation in the observation of exotic cascade hyperons can be found in the work by H G Fischer and S Wenig, hep-ex/0401014.

(2) In addition to the component of the antidecuplet $\Xi^*_{3/2}$ with mass 1.78 – 1.79 GeV, indicated in Table 2 and in Fig. 2, chiral soliton models also predict a state with the same isospin $(I = 3/2)$ and strangeness $(S = -2)$ that belongs to the {27}-plet. The mass of this state, according to Ref. [12], is 1.85 GeV, which is very close to the values obtained by the NA49 collaboration [6].

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