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Baryons in the Nambu – Jona-Lasinio Model: A Review[1]

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Abstract

The Nambu - Jona-Lasinio model in its SU(2) and SU(3) versions with scalar and pseudoscalar coupling are applied to baryons. The parameters of the model are fixed in the meson sector. The baryons arise as a soliton of three valence quarks coupled to the Dirac sea. Within the SU(2) version the nucleon static properties including g_A, magnetic moments and nucleon electric polarizability as well as the electromaggnetic and axial form factors are described quite successfully. The nucleon-delta splitting comes out reasonable well. In SU(3) version the strangness carrying baryons are described as SU(3)-rotational excitations of the SU(2)-soliton embedded in the SU(3)-sector. The mass splittings between the octet and decuplet as well as within the multiplets are predicted in a good agreement with experimantal values. Isospin splittings and $g_A^{(0)}$, $g_A^{(3)}$ and $g_A^{(8)}$ are reproduced within the experimantal errors. In our calculations a higher order $(1/N_c)$ rotational corrections are taken into account.

In the last years because of the tremenduos difficulties in the nonperturbative (strong coupling) regime of QCD there is an increasing interest in the effective chiral models respecting the chiral symmetry breaking. Among these models the Nambu – Jona-Lasinio model applied to quarks is the only one which allows to combine in a

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natural way the two extreme opposite pictures - the pure valence one advocated by the constituent quark model and that used in the Skyrme model. The model makes use of the phenomenologicaly supported "standard" picture of the nucleon as a bound state of three valence quarks coupled to the mesons arising as quark-antiquarks pairs from the Dirac sea. Similar to the other effective models the NJL model [1] posseses two shortcomings. On the one hand in the model the quark confinement is not implemented and one can only hope that the confinement is not of great importanse for the very low-energy baryon properties. On the other hand the model is not renormalizable. It means that a finite cut-off is needed to make the theory finite. The cut-off introduces into the theory a physical scale which must be fixed by reproducing some physical quantity.

In this review we shortly present the formalism and some of the relevent results.

1. NJL model with SU(2)-flavour

We start with the simplest lagrangean of the NJL model with scalar and pseudoscalar quark-quark couplings [1]:

$$\mathcal{L} = \overline{\Psi} i\gamma^\mu \partial_\mu \Psi - m_0 \overline{\Psi}\Psi + \frac{G}{2}[(\overline{\Psi}\Psi)^2 + (\overline{\Psi}\gamma_5 \vec{\tau}\Psi)^2]. \tag{1}$$

Here Ψ describes a quark field with $SU(2)$-flavour (u and d) quarks, $N_c = 3$ colours and the average current quark mass $m_0 = (m_u + m_d)/2$. Introducing auxiliary sigma and pion fields by $\sigma = -g\overline{\Psi}\Psi/\lambda^2$ and $\vec{\pi} = -g\overline{\Psi} i\gamma_5 \vec{\tau}\Psi/\lambda^2$ and assuming them to be classical (zero boson loop approximation), one can use path integral techniques in euclidean metric to evaluate the corresponding effective action as a sum of a quark part

$$S^q_{eff} = Sp\log\left(\frac{\partial}{\partial t} + h - \mu\right) \qquad \text{with} \qquad h = \frac{\vec{\alpha}.\vec{\nabla}}{i} + \beta g(\sigma + i\gamma_5 \vec{\pi}.\vec{\tau}) \tag{2}$$

with μ being the chemical potential, and a meson one

$$S^m_{eff} = \int d^4x \left\{\frac{\lambda^2}{2}(\sigma^2 + \vec{\pi}^2) + m_\pi^2 f_\pi \sigma\right\}. \tag{3}$$

In the meson part using the PCAC the current mass m_0 is eliminated in favor of the pion mass m_π and the pion decay constant f_π. The coupling constant G is related the new one λ by $G = g^2/\lambda^2$ where the additional coupling constant g is introduced for convinience. The latter will be fixed later to the physical coupling constant in order to identify the auxiliary pion and sigma fields with the physical ones in the meson sector. Only one part of the effective action, $S^q_{eff}(\mu = 0)$, coming from the Dirac sea is divergent. Using the proper time scheme [2] we make it finite. Subtracting the vacuum contribution one can express the energy (see for more details refs. [3] as a sum of a quark part

$$E^q = \sum_{0\le\epsilon_\alpha\le\mu} \epsilon_\alpha + \frac{N_c}{2}\left\{\sum_\alpha R_{3/2}(\epsilon_\alpha, \Lambda) - \sum_\alpha R_{3/2}(\epsilon^0_\alpha, \Lambda)\right\} \tag{4}$$

with the proper-time regularization function

$$R_\alpha(\epsilon, \Lambda) = \frac{1}{\sqrt{4\pi}} \int_{1/\Lambda^2}^{\infty} \frac{d\tau}{\tau^\alpha} \exp(-\epsilon^2\tau), \tag{5}$$

and the meson part which up to a trivial factor is given by eq.(3). The energies ϵ_α are the eigenvalues of the hamiltonian (2) and ϵ^0_α correspond to the vacuum solution.

We fix the parmeters of the model, namely the coupling constant λ^2 and the cut-off Λ, in the meson sector of the vacuum by reproducing the experimental values of the pion decay constant f_π and the pion mass m_π. This procedure originates from Eguchi [4] and is described in detail in ref. [3]. We use the stationary meson field configuration $< \vec{\pi} >_v = \vec{\pi}_v = 0$ and $< \sigma >_v \neq 0$ where the the latter is given by the non-trivial solution of the well-known gap equation. It generates the constituent quark mass $M = g < \sigma >_v = gf_\pi$. Actually we use the gap equation to express the coupling constant λ^2 as a function of the constituent mass M. For the meson dynamics one can use the meson propagators [5, 6] and to fix the g to the physical pion coupling constant in order to identify the auxiliary σ and $\vec{\pi}$ with the physical ones. The next step is to reproduce the physical value of the pion decay constant. The latter leads to a relation between the M and Λ. All those conditions together leave the vacuum value of the constituent mass M as the only free parameter. In principle, the empirical values of the quark condensate and the quark bare mass can be used to fix M but actually they still leave a broad range for M.

For the solitonic sector we solve the Dirac equation (3) with the equations of motion of the meson fields:

$$\sigma = \frac{g}{\lambda^2} N_c \left\{\sum_\alpha \overline{\phi}_\alpha \phi_\alpha R_{1/2}(\epsilon_\alpha, \Lambda) - \sum_{0\le\epsilon_\alpha\le\mu} \overline{\phi}_\alpha \phi_\alpha\right\} - \frac{f_\pi m_\pi^2}{\lambda^2}. \tag{6}$$

$$\pi = \frac{g}{\lambda^2} N_c \left\{\sum_\alpha \overline{\phi}_\alpha i\gamma_5 (\vec{\tau}.\hat{r}) \phi_\alpha R_{1/2}(\epsilon_\alpha, \Lambda) - \sum_{0\le\epsilon_\alpha\le\mu} \overline{\phi}_\alpha i\gamma_5 (\vec{\tau}.\hat{r}) \phi_\alpha\right\}. \tag{7}$$

The meson fields are assumed to be in a hedgehog form and are restricted on the chiral circle $\sigma^2 + \vec{\pi}^2 = f_\pi^2$. We use a numerical self-consistent iterative procedure [3] based on a method proposed by Ripka and Kahana [7].

As a next step the classical soliton should be quantized. To that end one can make use of the rotational symmetry of the solution introducing isorotating hedgehog meson fields

$$\tilde{U}(x) \to R(t) U(x) R^\dagger(t), \tag{8}$$

where $R(t)$ is a unitary SU(2) matrix. Following the usual canonical quantization procedure [8] one can relate the angular velocities to the spin (isospin) operators by

$$\theta\Omega^i(\Omega^a) \to -iI^i(T^a). \tag{9}$$

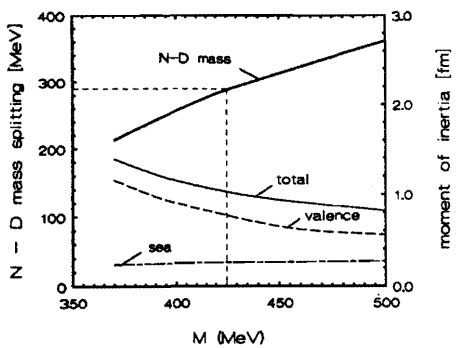

Figure 1: Nucleon-Delta mass splitting and the moment of inertia as functions of the constituent quark mass.

Thus one can assign proper spin and isospin quantum numbers to the soliton and evaluate the nucleon properties like energies, radii, form factors, spin and isospin contents, etc.

For the energy an expansion up to the first non-vanishing order with respect to angular velocity gives the well-known rotational energy correction to the classical energy [9]

$$E^I = E_{sol} + \frac{I(I+1)}{2\theta} - \frac{< T^2 >}{2\theta} - \frac{< P^2 >}{2\theta}. \tag{10}$$

One should keep in mind that from this energy the corresponding translational and rotational spurious zero-point energies $\frac{<T^2>}{2\theta} + \frac{<P^2>}{2\theta}$ have to be subtracted. Here θ is the moment of inertia given as a sum of valence and sea part [9]. The latter is reguralized in the proper time scheme. The results of the mass $N - D$ splitting and the moment of inertia are shown in 1.

The results concerning the nucleon static properties are listed in 1 [10, 11]. As can be seen there is an overall good agreement with the experimantal numbers for M around 420 MeV.

Some recent calculations [12] (theese Proceedings) are done for the electric polarizability of the nucleon. Rotational $1/N_c$ corrections, analogous to those, discussed for the case of g_A, heve been also studied and found to be important – their contribution is about 30% of the total α. Since the nucleon polarizability is strongly correlated to g_A, it is plausible that including the rotational corrections the model gives value of g_A [11] in good agreement with experiment. The Dirac sea effects, describing the physics of the pion cloud, are found to be dominating for α in qualitative agreement with other models. After appriximately including additional corrections due to the

Table 1: Nucleon static properties for different constituent quark masses.

Quantity	Constituent Quark Mass 370 MeV		420 MeV		450 MeV		Experiment
	total	sea	total	sea	total	sea	
$< r^2 >_{T=0}$ $[fm^2]$	0.63	0.05	0.52	0.07	0.48	0.09	0.62
$< r^2 >_{T=1}$ $[fm^2]$	1.07	0.33	0.89	0.41	0.84	0.45	0.86
$< r^2 >_n$ $[fm^2]$	–0.22	–0.14	–0.18	–0.17	–0.18	–0.18	–0.12
$< r^2 >_p$ $[fm^2]$	0.85	0.19	0.70	0.24	0.66	0.27	0.74
$\mu_{T=0}$ $[n.m.]$	0.68	0.09	0.62	0.03	0.59	0.05	0.88
$\mu_{T=1}$ $[n.m.]$	3.74	0.96	3.53	1.07	3.47	1.17	4.71
μ_p $[n.m.]$	2.21	0.53	2.08	0.55	2.03	0.61	2.79
μ_n $[n.m.]$	–1.53	–0.44	–1.46	–0.52	–1.44	–0.56	–1.91
$M_\Delta - M_N [MeV]$	213	—	280	—	314	—	294
g_A	1.32	0.15	1.29	0.20	1.29	0.23	1.26

N-Δ mass splitting we obtain for the total polarizability $\alpha \simeq 19 \times 10^{-4}$ fm^3, a value which is about a factor of two larger than the experimental one. This discrepancy, although large, is rather small compared to all other chiral model calculations providing reasonable values for g_A as well.

The nucleon electromagnetic form factors are given by

$$G_E(q^2) = \langle N(p) | j_o^{em}(0) | N(p') \rangle$$
$$\frac{1}{2M_N} G_M(q^2) \varepsilon_{ikl} q^k \langle N | \sigma^l | N \rangle = \langle N(p) | j_i^{em}(0) | N(p') \rangle, \tag{11}$$

where $| N(p) \rangle$ is the quantized nucleon state and j_o^{em} is zero-component of the electromagnetic current operator defined as

$$j_\mu^{em} = \frac{\delta}{A_\mu} S_{eff}^f [A_\mu, \Omega]. \tag{12}$$

To that end an external electromagnetic field A_μ is coupled. Similar to the moment of inertia the form factors incude also a valence and sea contribution. It should be noted that the isoscalar form factor does not need any regularization which is not the case of the isovector one. The proton and neutron electromagnetic form factors calculated [10] are shown in Fig.2, Fig.3, Fig.4, and Fig.5, respectively. The proton electromagnetic form factors as well as the neutron magnetic one are in good agreement with the experiment for $M \approx 420$ MeV whereas the neutron electric form

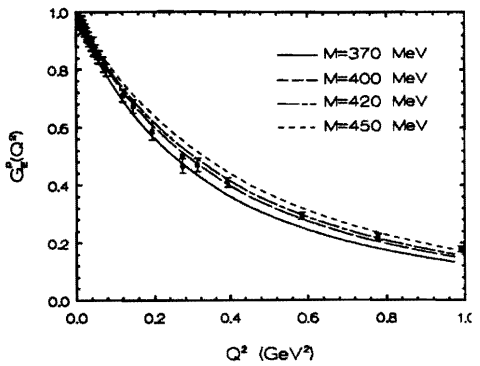

Figure 2: Proton electric form factor as a function of the constituent quark mass.

factor is overestimated. For the proton form factors the sea contribution is less than 5% which is not the case of the neutron electric form factor where the valence quarks and the sea contribute with similar magnitude and opposite signs.

2. NJL model with SU(3)-flavour

The SU(3)-invariant NJl model withscalar and pseudoscalar coupling reads

$$\mathcal{L}_{NJL} = \bar{q}(x)(i\partial\!\!\!/ - m)q(x) - \frac{G}{2}\left[(\bar{q}(x)\lambda^a q(x))^2 + (\bar{q}(x)i\gamma_5\lambda^a q(x))^2\right] \qquad (13)$$

where $m = \mathrm{diag}(m_u, m_d, m_s) = m_1\mathbf{1} + m_3\lambda_8$ is the quark mass matrix and λ_a are the usual Gell-Mann matrices with $\lambda^0 = \sqrt{(\frac{2}{3})}\mathbf{1}$. Actually we use the current quark masses $m_0 = \frac{1}{2}(m_u + m_d)$. Integrating out the quarks the effective action is given by

$$S_{eff} = -\mathrm{Sp}\ \log\left(i\partial\!\!\!/ - m - g(\sigma^a\lambda^a + i\gamma_5\pi^a\lambda^a)\right) + \frac{\mu^2}{2}(\sigma^a\sigma^a + \pi^a\pi^a) \qquad (14)$$

Hereafter we work in the chiral limit and treat the strange quarkmass m_s perturbatively in first order. Similar to the SU(2) case we fix the parameters in the meson sector of the vacuum by reproducing the experimental values of the meson masses $m_\pi = 139MeV$ and $m_K = 496MeV$, as well as the pion decay constant $f_\pi = 93MeV$. This results in the relation $m_K^2/m_\pi^2 = (m_s + m_0)/2m_0$ with $m_0 = \frac{1}{2}(m_u + m_d)$. In order to fulfill it one needs a current mass m_0 of about 6 MeV. To this end we use a proper time regularization with a modified weight function. The latter leads to a higher value of the Σ-commutator $\Sigma = 51$ MeV.

For the SU(3) soliton we use the $SU(2)$-one U_0 trivially embedded in SU(3)

$$U_2 = \begin{pmatrix} U_0 & 0 \\ 0 & 1 \end{pmatrix}$$

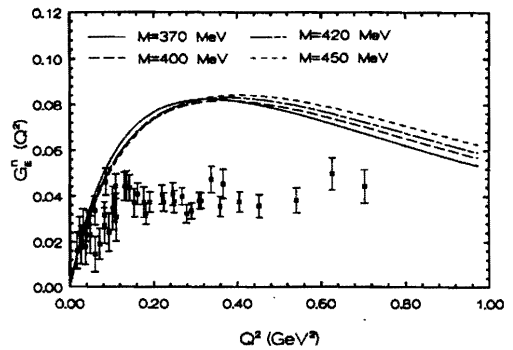

Figure 3: Neutron electric form factor as a function of the constituent quark mass.

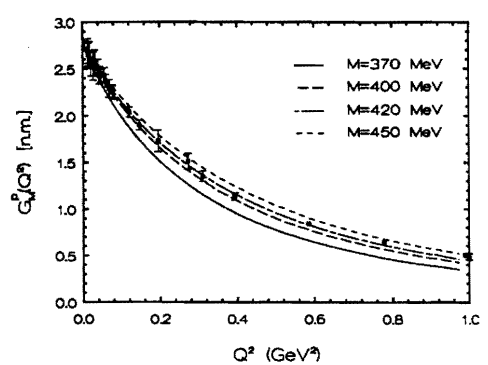

Figure 4: Proton magnetic form factor as a function of the constituent quark mass.

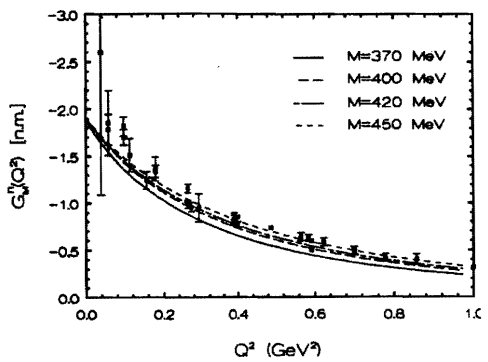

Figure 5: Neutron magnetic form factor as a function of the constituent quark mass.

In order to quantize the classical soliton we follow the same scheme (canonical quantization procedure) as in the SU(2) case (see eq.(8)) but with a SU(3) isorotating meson fields. As a next step we expand the effective action in Ω_A up to second order

$$L^{rot} = \frac{1}{2} I_{AB}\Omega_A\Omega_B - \frac{N_c}{2\sqrt{3}}\Omega_8, \tag{15}$$

where I_{AB} is the $SU(3)$ tensor of the moment of inertia, as well as up to the first order in m_s

$$L^{(1)} = -\frac{2}{3}\frac{m_s}{m_u + m_d}\Sigma\left(1 - D_{88}^{(8)}(R)\right) \tag{16}$$

and $m_s\Omega_A$

$$L^{(2)} = -\frac{2m_s}{\sqrt{3}} K_{AB} D_{8A}^{(8)}(R)\Omega_B. \tag{17}$$

All other corrections are neglected. Here $D_{8A}^{(8)}(R)$ is the SU(3)-Wigner function and by K_{AB} we denote the tensor of the anomalous moments of inertia. These moments need no regularization, because they originate from the imaginary part of the effective Euclidean action and so they are finite. It should be stressed that the appearance of the linear term in Ω_8 in (15) is due to the discrete valence level in the spectrum whereas in the Skyrme model a term like this can be obtained only by adding the Wess-Zumino term.

Now we quantize canonically by defining $SU(3)$ right generators $R_{A=1,\ldots,8}$ by $R_A = -\partial L/\partial\Omega_A$, where $L = L^{rot} + L^{(1)} + L^{(2)}$. Using the collective hamiltonian, obtained in approximation $\mathcal{O}(m_s^2)$, in a strict perturbation theory up to second order in m_s we calculated [13] the masses of the member of the octet and decouplet. We fixed m_{Σ^*} to the experimental value, because as in the $SU(2)$-case the total masses are always

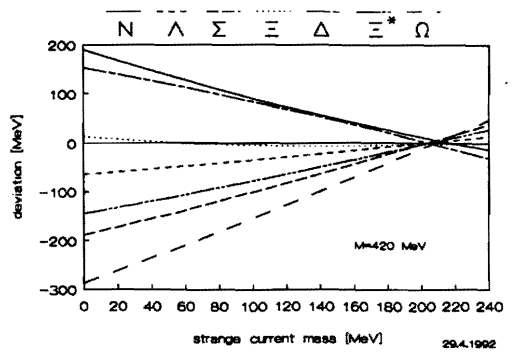

Figure 6: Deviations of the theoretical predicted masses from experimental vlaues as a function of the strangness quark mass at constituent quark mass $M = 420$ MeV.

too large by a constant shift (zero-point energies are not subtracted). It should be stressed that we obtained the Gell-Mann Okubo relations

$$2(m_N + m_\Xi) = 3m_\Lambda + m_\Sigma \qquad \text{and} \qquad m_\Omega - m_{\Xi^*} = m_{\Xi^*} - m_{\Sigma^*} = m_{\Sigma^*} - m_\Delta \tag{18}$$

as well as the Guadignini formula [16]

$$m_{\Xi^*} - m_{\Sigma^*} + m_N = (1/8)(11m_\Lambda - 3m_\Sigma) \tag{19}$$

without any additional assumptions. The latter is derived in the Skyrme model only by adding a term associated to the hypercharge simply by hand. Our results are illustrated in 6 where the discrepancy between the theory and experiment for the masses of the members of the octet and decouplet is presented as a function of the strangness quark mass. The hadronic parts of the their isospin mass differences are reproduced [14] quite successfully. The spin-structure of the nucleon reveals itself in the our estimates [15] $g_A^{(0)} = 0.37$, $g_A^{(3)} = 1.37$ and $g_A^{(8)} = 0.29$ compared to the experimental values $g_A^{(0)} = 0.24 \pm 0.09$, $g_A^{(3)} = 1.25$ and $g_A^{(8)} = 0.35 \pm 0.04$.

3. Summary

The numerical results for the solitonic sector of the NJL model are basically independent on the regularization scheme used as long as the cut-off is treated as a parameter of the system, which is fixed in the meson sector. The calculations of the literature, all being based on hedgehog structure and physical values for the meson sector, can be summarized as follows (status autumn 1993):

i) SU(2), σ and π field, chiral circle: For a constituent quark mass $M = 420$MeV the following observables within 15% are reproduced: isoscalar and isovector

charge squared radius of the nucleon, $< r^2 >^c_{T=0}$ and $< r^2 >^c_{T=1}$, nucleon sigma term, Σ, nucleon axial coupling constant and magnetic moments of proton and neutron (if the recent $1/N_c$ corrections are included), g_A, μ_p and μ_n, the nucleon-delta splitting, $M_\Delta - M_N$, the q-dependence of the form factor $G^p_E(q^2)$, $G^p_M(q^2)$, $G^n_M(q^2)$, $g_A(q^2)$. The nucleon energy comes out at about 800MeV if rotational and translational zero modes corrections are included. The neutron squared radius, being a very sensitive quantity, is by 50% too large than the experimental value, the $G^n_E(q^2)$ is also overestimated at finite q^2.

ii) SU(3), σ, π, K and η fields. chiral circle for sigma and pion, trivial embedding of SU(2) into SU(3), perturbative treatment of m_s up to second order or Yabu-Ando approach: Splitting between and within spin 1/2 and spin 3/2 baryons is reproduced within few MeV. The splitting within isospin multiplets is reproduced and a common value for $m_u - m_d$ is found. If the zero point corrections to translation and the SU(3)-rotations are included the energies of all octet and decuplet baryons are too large by a constant shift of approximately 100 MeV. If recent $1/N_c$ corrections are included the g^0_A, g^3_A and g^8_A of the nucleon are reproduced within within experimental errors. The nucleon Σ-term is reproduced with a strangeness content of 15% and the recently measured Gottfried sum comes out as well.

iii) SU(2), σ, π, ρ, A_1 and ω mesons, chiral circle for sigma and pion. None of the present approaches has been developped to an extent that one can judge the influences or the necessity of vector mesonic coulings. The observables calculated by now are the isoscalar charge radius of the nucleon and g_A, both deviating noticebly from experiment. There seems to be also conceptual problems with treating the omega meson. On the other hand the repulsive character of the omega guarantee a stable soliton and one does not require any more the restriction to the chiral circle.

The problem encountered by now can be summarized as following: The semiclassical quantization procedure is not perfectly understood yet, although it is common usage. Its relationship to well-established quantum-mechanical theories like Peierls-Yoccoz projection approaches is yet unknown. If one looks at g_A and magnetic moments of the nucleon its convergence in the rotational angular velocity Ω seems to be slow. Furthermore the PCAC relation is slightly violated within the semiclassical quantization scheme when the rotational corrections, linear order of Ω, are included.

The quantum corrections due to zero modes of translation and rotation add up to $30 - 40\%$ of the resulting baryon mass and hence by far too large for a correction. The corresponding corrections of non-zero modes are of order $\mathcal{O}(N_c)$ as well and they are not known yet. There are indications in the Skyrme model that they are small, but a clear statement in the NJL model is missing.

Altogether one can summarize: Quark models like the NJL model with a polarized Dirac sea and involving Goldstone type couplings (i.e. non-dynamic mesons), have

been by now quite successful in describing the ground states of baryons in the octet and decuplet.

If one ignores vector mesons, whose relevance in the present model is still under debate, the calculations support a clear picture of the baryons consisting of three localized valence quarks interacting with a moderately polarized Dirac sea. This is in contrast to Skyrme-like models, which are governed by meson fields, exhibiting certain topological properties, and quark-meson models, where valence quarks are coupled to dynamical meson fields.

The problem encountered concern the semiclassical quantization method. It has certain merits, but it deserves improvement. Last but not least, the NJL-model lacks confinement which should be necessary for radially excited states of baryons.

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