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**Nucleon-Antinucleon Annihilation
in Chiral Soliton Model**

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NUCLEON-ANTINUCLEON ANNIHILATION IN CHIRAL SOLITON MODEL

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Abstract

We investigate annihilation process of nucleons in the chiral soliton model by the path integral method. A soliton - antisoliton pair is shown to decay into mesons at range of about 1fm, defined by the $S\bar{S}$ potential. Contribution of the annihilation channel to the elastic scattering is discussed.

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1. Introduction

The understanding of nucleon-antinucleon annihilation processes is an important issue of nuclear physics (for the recent analysis of experimental data see [1]). The traditional way of describing the nucleon-antinucleon interactions at low energies is the potential methods. Purely real part of the potential is usually explained by the meson exchange and imaginary part is accounted phenomenologically. Another semiphenomenological approach is based on the coupled channels method. These methods give the value of the annihilation radius of about $0.8 \div 1.2$ fm [8],[10].

Applications of chromodynamics (QCD), which is believed to be a theory of strong interaction at high energy, in this case obviously require models. Among them were discussed an effective perturbative approach in the flux-tube type quark models [2], [3] and the nonperturbative one in the chiral soliton model [4].

Chiral models like the Skyrme model [5] seem to be more appropriate for description of low - energy hadron physics. They represent a theory for interacting pion field which has the topological soliton solutions. The chiral theory may be considered as being a low - energy limit of the QCD. The chiral Lagrangian contains only physical degrees of freedom (meson fields) and realizes fundamental symmetries of the initial QCD Lagrangian. In particular, it reproduces the phenomena of the spontaneous breaking of chiral symmetry. The connection of the chiral theory with QCD was proved by calculation in the $N_c \rightarrow \infty$ limit [6].

In paper [7] an analogous model, derived from the low - energy effective action of instanton model of the QCD vacuum, was proposed.

In the chiral theory nucleons are identified with the low - lying rotating states of the topological soliton (Skyrmion).

When a nucleon and an antinucleon come close to each other, they interact by means of their external pion clouds. Of course, because of the pion exchange, the nucleon states can transform into other higher rotational states (for example, Δ - isobar). If the distance between the centres of nucleon and antinucleon becomes less than some critical radius, this state looks like an excited state of the pion field over the perturbative vacuum rather than the nucleon - antinucleon pair. Then this state decays into pions.

So, an important parameter of the theory is the critical radius R_0 , which is to be identified with the radius of annihilation. At large distances ($R > R_0$) solitons can be described by the collective coordinates - the coordinates of the centres and the parameters of orientation. At small distances ($R < R_0$) the collective coordinates lose their meaning and the total pion field of the resulting configuration must be considered as being a fluctuation over trivial vacuum. The critical radius R_0 will be shown to be determined from the form of classical soliton - antisoliton potential and appears to be $R_0 \sim 1$ fm.

On the other hand, the phenomenology of the low - energy nucleon - antinucleon interactions gives the value of the radius of annihilation of about $0.8 \div 1.2$ fm[8, 9, 10].

A natural way of describing the process of soliton - antisoliton annihilation is the path integral method. Using it we construct in this paper the low - energy amplitude of $N\bar{N}$ - annihilation and obtain the expression for contribution of the annihilation channel to the amplitude of low - energy elastic scattering. The crucial point of our discussion is the necessity of redefining the semiclassical quantization procedure at small distances.

The paper is organized as follows. In section 2 we reproduce the main formulae of the chiral soliton model. In section 3 we give some arguments for redefinition of collective integrating measure at small distances. In section 4, it is shown how the annihilation amplitude can be evaluated.

2. General formulae

We will consider the $SU(2)$ - chiral theory.

The Lagrangian of the chiral model has the form of the nonlinear σ - model describing isovector pion fields $\vec{\pi}$:

$$L = \frac{1}{4} f_\pi^2 \text{Tr} L_\mu L^\mu + \frac{1}{32e^2} \text{Tr}[L_\mu, L_\nu]^2, \quad (1)$$

where $L_\mu = U^\dagger \partial_\mu U$ and $SU(2)$ - valued matrix U can be represented as $U = e^{i\vec{\tau}\vec{\Phi}} = \frac{1}{f_\pi}(\sigma + i\vec{\tau}\vec{\pi})$. The soliton solution with unit baryon number is obtained with the hedgehog ansatz [5]:

$$U_{N,0}(\vec{x}) = \exp(i\vec{\tau} \cdot \frac{\vec{x}}{|\vec{x}|} \theta(|\vec{x}|)), \quad (2)$$

where θ is a scalar function. We use the following determination of the antiskyrmion solution: $U_N = U_N^\dagger$.

The equations of motion stay invariant under translation of the centre of the soliton and under rotation produced with multiplying U by any constant $SU(2)$ - matrices A : $U = A^\dagger U_0 A$.

3. Quasi-zero mode integration

Let us consider quantization of the static soliton - antisoliton configuration. We start from the configuration corresponding to the well separated soliton- antisoliton pair using "production ansatz":

$$U(\vec{x}; \vec{R}_N, \vec{R}_{\bar{N}}, \vec{\omega}_N, \vec{\omega}_{\bar{N}}) = U_N(\vec{x} - \vec{R}_N, \vec{\omega}_N) \cdot U_{\bar{N}}(\vec{x} - \vec{R}_{\bar{N}}, \vec{\omega}_{\bar{N}}) = e^{i\vec{\tau}\vec{\Phi}_{N\bar{N}}(\vec{x})}, \quad (3)$$

$$U_N(\vec{x}, \vec{\omega}) = A^\dagger(\vec{\omega}_N) \cdot U_{N,0}(\vec{x}) \cdot A(\vec{\omega}_N),$$

where $A(\vec{\omega}) = e^{i\vec{\tau}\vec{\omega}}$ is the matrix of the soliton orientation. In the $|\vec{R}_N - \vec{R}_{\bar{N}}| \rightarrow \infty$ limit, the configuration of the field $\vec{\Phi}_{N\bar{N}}$ is the solution of the equation of motion with energy $E = 2M_N$ and realizes the local minimum of the action. Note that the absolute minimum in zero topological sector is achieved on the trivial configuration $\vec{\Phi} = 0$.

When one quantizes fluctuations of the pion field around the classical field (3) $\vec{\phi}(\vec{x}, t) = \vec{\Phi}(\vec{x}, t) - \vec{\Phi}_{N\bar{N}}(\vec{x})$, $\vec{\pi} = f_\pi \vec{\phi}$ in the semiclassical approximation he faces the well known problem of existence of flat directions (valleys) in the functional space. It means that there are zero frequencies in the spectrum of the operator $\delta^2 S / \delta \vec{\Phi}^2$. Because of this fact, the integration along these directions are not Gaussian and must be carried out exactly. A usual way to treat zero modes is to rewrite these integrals (using the Faddeev - Popov technique) as a path integral over collective coordinates $\vec{R}_i(t)$, $\vec{\omega}_i(t)$ [11]

$$\int D\vec{\pi} e^{iS(\vec{\Phi}_{N\bar{N}}, \vec{\pi})} = \int \prod_{i=N, \bar{N}} D\vec{R}_i(t) D\vec{\omega}_i(t) \int D\vec{\pi}' e^{iS(\vec{\Phi}_{N\bar{N}}, \vec{\pi}')}. \quad (4)$$

Boundary conditions corresponding to a soliton - antisoliton pair are implied. Prime in the right-hand side denotes that the integration is carried out over the functional subspace of nonzero modes. We interest in zero modes only, regarding nonzero ones as giving rise just to renormalization of the soliton mass.

Factorization of the integration measure over the collective coordinates of the well-separated soliton- antisoliton pair onto product of measures of individual solitons $d\mu = d\mu_N \cdot d\mu_{\bar{N}}$ in Eq.(4) is caused by the fact that $S(\vec{\Phi}_{N\bar{N}})$ is invariant under independent translations of the centres and under rotations of a soliton or an antisoliton. When the static field configuration describes the

soliton and antisoliton placed at large but finite distance it does not satisfy the equation of motion because of its nonlinearity. The corresponding action now is the function of the relative distance $\vec{R} = \vec{R}_N - \vec{R}_{\bar{N}}$ and of the matrix of relative orientation $A = A_N^+(\vec{\omega}_N) \cdot A_{\bar{N}}(\vec{\omega}_{\bar{N}}) = e^{i\vec{\tau}\vec{\omega}} = a_0 + i\vec{\tau}\vec{a}$:

$$S_{N\bar{N}} = S_N + S_{\bar{N}} - \int V_{N\bar{N}}(\vec{R}, \vec{\omega}) dt. \quad (5)$$

Eq.(5) defines the soliton - antisoliton potential $V_{N\bar{N}}$. It was evaluated numerically in Ref.[12] as a function of the relative distance and of two orientation parameters $\vec{a}^2$ and $T = \frac{\vec{a} \cdot \vec{R}}{|\vec{a}| \cdot |\vec{R}|}$ (See Fig 1).

We see that now some valleys in the functional space are not flat. However, at sufficiently large R , the variational derivative $\delta^2 S / \delta \Phi^2$ taken at variation of $\vec{\Phi}$ along the valley is much less than that along the transverse direction. Therefore, quasi-zero modes can be treated in the same way as zero ones [13], with account of the potential $V_{N\bar{N}}$. We conclude that the measure of integration over all collective coordinates corresponding to both zero and nonzero modes at large R reads:

$$\int \prod_{i=N, \bar{N}} D\vec{R}_i(t) D\vec{\omega}_i(t) \exp \left\{ i \int dt \left(-2M_N + \frac{M_N}{2} (\dot{\vec{R}}_N^2 + \dot{\vec{R}}_{\bar{N}}^2) + \frac{I_N}{2} (\dot{\vec{\omega}}_N^2 + \dot{\vec{\omega}}_{\bar{N}}^2) - V_{N\bar{N}}(\vec{R}, \vec{\omega}) \right) \right\}. \quad (6)$$

Evidently, it has the form of a path integral describing scattering of two quantum particles - nucleon and antinucleon interacting with each other by means of the potential $V_{N\bar{N}}$.

Expression (6) needs some redefinition at small distances between the centres of a soliton and an antisoliton. The main observation is that the static field configuration (3) at the limit $R \rightarrow 0$ is not distinguished from a meson field fluctuation around the trivial vacuum $\vec{\Phi} = 0$. In fact, since there is strong attraction in the system when R becomes less than the radius of soliton R_0 [12] (See Fig.1), the quantization procedure breaks due to loss of stability along the quasi-zero mode direction. This configuration is far from being a solution of the equation of motion and cannot be used as a stationary phase point in the perturbation theory expansion. The collective coordinate method (and, generally speaking, soliton - antisoliton picture) in this situation is not adequate. Moreover, when soliton and antisoliton come close, regions with opposite topological charges are nearly overlapping. Therefore, a region in functional space corresponding to the soliton - antisoliton configurations at $R < R_0$ is to be regarded as perturbative quantum fluctuations over trivial vacuum $\vec{\Phi} = 0$. However, these fluctuations are already taken into account in the path integral analogous to left-hand side of Eq. (3) but with zero topological boundary conditions. To avoid double count, this region must be excluded from Eq. (6).

To formalize this statement, let us represent the path integral measure over $\vec{R}_N(t), \vec{R}_{\bar{N}}(t)$ in the Eq.(6) in the following form (in accordance with its definition):

$$D\vec{R}_N(t) D\vec{R}_{\bar{N}}(t) = \prod_t \left(\frac{d^3 \vec{R}_N(t) d^3 \vec{R}_{\bar{N}}(t)}{(2\pi)^3} \right). \quad (7)$$

In order to exclude the short range region from the integration over $\vec{R}_i(t)$, we put (in the first approximation)

$$d^3 \vec{R}_N(t) d^3 \vec{R}_{\bar{N}}(t) \rightarrow d^3 \vec{R}_N(t) d^3 \vec{R}_{\bar{N}}(t) \cdot \Theta(R(t) - R_0). \quad (8)$$

for every moment of the time t (ansatz proposed by Jevicki [14]). Here $\Theta(x)$ is the standard Θ - function. Note that we neglect the dependence of the potential on the orientation parameters and use averaging value of R_0 . Eq.(8) means that the wave function of the nucleon - antinucleon pair vanish with its derivatives inside the sphere $|\vec{R}| = R_0$. This statement satisfactorily agrees with phenomenological data, provided that we put $R_0 \sim 0.8 \div 1.2 \text{ fm}$.

Assembling Eqs.(6) and (8) we see that the contribution of the annihilation channel to the scattering amplitude is provided by the trajectories coming from $R(t = -\infty) = \infty$ to the sphere $R(t_0) = R_0$ at some moment t_0 and outgoing from the sphere $R(t'_0) = R_0$ ($t_0 > t'_0$) to $R(+\infty) = \infty$.

4. Annihilation

Let us come to evaluating the annihilation amplitude. We shall consider the amplitude of soliton - antisoliton annihilation to vacuum in the presence of external source:

$$Z(J, \Psi_{N\bar{N}}) = \int D\tilde{\Phi} e^{i(S(\Phi) + \int d^3\vec{x} dt \tilde{\Phi}(\vec{x}, t) J(\vec{x}, t))} \cdot \Psi_{N\bar{N}}^{in}. \quad (9)$$

$\Psi_{N\bar{N}}^{in} = e^{i(P_N X + P_{\bar{N}} Y)} \cdot \mathcal{D}(\omega_N) \mathcal{D}(\omega_{\bar{N}})$ is the in-state wave function of a nucleon - antinucleon pair, where $\mathcal{D}$ is the wave function of a soliton rotating state. Note that $\mathcal{D}$ is fixed by the spin - isospin quantum numbers of a baryon.

It is implied that integration in Eq. (9) is carried out over fields which satisfy boundary conditions corresponding to the presence of a soliton - antisoliton pair at $t = -\infty$ and pion radiation condition at $t = +\infty$. Process of annihilation proceeds through two stages. At the first one a soliton and an antisoliton move from in-state ($R = \infty$) and at a moment t_0 turn up at a critical distance $R(t_0) = R_0$. Then, at the second stage, the field configuration $\Phi(\vec{x}, t) = \Phi(\vec{x}; \vec{R}_i(t_0), \vec{\omega}_i(t_0))$ (See Eq. (3)) decays into pions. Summing up amplitudes with different t_0 and extracting variables related to overall orientation we obtain the amplitude of $N\bar{N}$ annihilation into n pions (in CM frame):

$$\int d\Omega \delta(|\vec{R}| - R_0) \langle n\pi | \tilde{\Phi}(\vec{x}, t_0; \vec{R}, \vec{\omega}) \rangle \langle \vec{R}, \vec{\omega}, t_0 | \Psi_{N\bar{N}}^{in} \rangle, \quad (10)$$

where $d\Omega$ denotes the integration measure over all variables: $d\Omega = \int d^4X d^4Y \int dt_0 d^3\vec{x} \int d^3\vec{R} d^3\vec{\omega}$.

The matrix element

$$\langle \vec{R}, \vec{\omega}, t_0 | \Psi_{N\bar{N}}^{in} \rangle = \langle \vec{R}, \vec{\omega} | e^{-iH_{int}(t_{in}-t_0)} | \Psi_{N\bar{N}}^{in} \rangle \quad (11)$$

has a meaning of value of the $N\bar{N}$ quantum-mechanical wave function at the point $\vec{R}, \vec{\omega}$ and at time t_0 . Of course, it contains admixture of $N\bar{\Delta}$, $\Delta\bar{N}$, $\Delta\bar{\Delta}$ - states. It can be evaluated, for example, by means of the solution of the Lippmann - Schwinger equation with the long-range $N\bar{N}$ potential defined in Eq. (5), and with boundary condition at $|\vec{R}| = R_0$ to be evaluated from short - range dynamics.

Now, we wish to compute the decay amplitude of the state $\tilde{\Phi}$. We can do this by using the method of oscillator quantization [15]. So, if interaction of pions in the final state is neglected, the field $\tilde{\Phi}$ can be represented as a set of quantum oscillators $\tilde{\Phi}_{\vec{k}}$ numbered by $\vec{k}$ with frequency $\omega_k = |\vec{k}|$, where $\tilde{\Phi}_{\vec{k}}$ is the Fourier component of the field

$$\tilde{\Phi}(\vec{x}, t) = \int \frac{d^3\vec{k}}{(2\pi)^{3/2}} \left(\tilde{\Phi}(\vec{k}) e^{-i(\vec{k}\vec{x} - \omega_k x_0)} + \tilde{\Phi}^+(\vec{k}) e^{+i(\vec{k}\vec{x} - \omega_k x_0)} \right). \quad (12)$$

We evaluate $\langle n\pi | \tilde{\Phi} \rangle$ with neglecting the final state interaction. We wish to compute the amplitude of decay of the definite classical field configuration to the state with the definite number of pions. We can do this by using the method of oscillator quantization [15].

Let us consider the system as placed in a large box. Then, integral in Eq.(13) transforms to the sum over modes with discrete $\vec{k}$. Path integral description for such a system is equivalent to the Shrödinger equation with $H = \sum H_{\vec{k}}$, $\Psi = \prod_{\vec{k}} \Psi_{\vec{k}}(\tilde{\Phi}_{\vec{k}}, t)$, $\Psi_{\vec{k}}(\tilde{\Phi}_{\vec{k}}, t) = \Psi_{\vec{k}}(\tilde{\Phi}_{\vec{k}}) e^{-i\omega_k t}$.

So, our task is reduced to the well-known problem of evaluating of the transition amplitude in $\vec{k}$ -th harmonic oscillator from the state with the definite "coordinate" $\vec{\Phi}_{\vec{k}}$ to the state with the definite number of quanta $n_{\vec{k}}$ (0 or 1). This amplitude is the value of the oscillator wave function $\Psi_{n_{\vec{k}}}(\vec{\Phi}_{\vec{k}})$. In the continuum limit, we obtain:

$$\langle n_{\vec{\pi}} | \vec{\Phi}(\vec{x}, t_0; \vec{R}, \vec{\omega}) \rangle = N \cdot \frac{1}{\sqrt{n!}} \cdot \prod_i^n \sqrt{2\omega_i} f_{\pi} \vec{\Phi}(\vec{k}_i) \cdot \exp \left(i \sum_i^n (\omega_i t_0 - \vec{k}_i \vec{x}_i) \right). \quad (13)$$

N is the normalization constant:

$$N = \exp \left\{ -f_{\pi}^2 \int \frac{d^3 \vec{k}}{(2\pi)^{3/2}} \omega_k \vec{\Phi}^+(\vec{k}) \vec{\Phi}(\vec{k}) \right\}. \quad (14)$$

Integrations over $d^3 \vec{x}$ and dt_0 yield δ^4 -function of momentum - energy conservation, and the integrations over the parameters of the total space and iso-space rotations (orientations) result in the spin - iso-spin conservation restrictions.

For example, in the case of isospin rotations, it can be derived as follows. The matrix of total orientation $B = e^{i\vec{I}\vec{\omega}_t}$ can be easily extracted from Eq.(3). Then, let us make rotation of the field configuration $\vec{\Phi}$: $|\vec{\Phi}(B, \dots)\rangle = e^{i\vec{I}\vec{\omega}_t} |\vec{\Phi}(B = 1, \dots)\rangle$. We obtain the matrix element $\langle n_{\pi} | \vec{\Phi} \rangle$ in the form:

$$\langle n_{\pi}; T, T_3 | \vec{\Phi}(B, \dots) \rangle = \langle n_{\pi}; T, T_3 | e^{i\vec{I}\vec{\omega}_t} | \vec{\Phi}(B, \dots) \rangle = \sum_{T_3''} D_{T_3, T_3''}^T(B) \langle n_{\pi}; T, T_3'' | \vec{\Phi}(B = 1, \dots) \rangle, \quad (15)$$

where $D_{T_3, T_3''}^T(B)$ is Wigner D -function. Wave function $\langle \vec{R}, \vec{\omega}_t, t_0 | \Psi_{N\bar{N}}^{in}; T, T_3' \rangle$ is treated in the same way. Integration over $\vec{\omega}_t$ gives immediately

$$\frac{1}{2T+1} \delta_{TT'} \delta_{T_3 T_3'} \langle n_{\pi}; T, T_3 | \vec{\Phi}(B = 1, \dots) \rangle \langle \vec{R}, \vec{\omega}_t = 0, t_0 | \Psi_{N\bar{N}}^{in}; T', T_3' \rangle. \quad (16)$$

Turning back to formula (8), we now able to write the short-range $N\bar{N}$ -potential. Repeating discussion of the end of Section 3, we immediately obtain contribution of the annihilation channel to the elastic scattering amplitude:

$$\int d\Omega d\Omega' \delta(|\vec{R}| - R_0) \delta(|\vec{R}'| - R_0) \cdot \langle \Psi_{N\bar{N}}^{out} | \vec{R}', \vec{\omega}', t_0' \rangle \times \\ \langle \vec{\Phi}(\vec{x}', t_0'; \vec{R}', \vec{\omega}') | \vec{\Phi}(\vec{x}, t_0; \vec{R}, \vec{\omega}) \rangle \cdot \langle \vec{R}, \vec{\omega}, t_0 | \Psi_{N\bar{N}}^{in} \rangle. \quad (17)$$

To evaluate remaining matrix element $\langle \vec{\Phi}(t_0') | \vec{\Phi}(t_0) \rangle$ one can also apply the method of Ref.[15].

This expression (via additional condition $\Psi_{N\bar{N}}(\vec{R})|_{|\vec{R}| < R_0} = 0$) defines the value of logarithmic derivative of the nucleon - antinucleon wave function $\Psi'_{N\bar{N}} / \Psi_{N\bar{N}}|_{|\vec{R}|=R_0}$.

Formulae (8), (13) and (17) are in accordance with the results obtained by using the coupled channel method in nonrelativistic quantum mechanics (See, for example, Ref.[16]). Namely, the Ψ -function of the $N\bar{N}$ -system can be decomposed into baryon and meson channels by means of the projecting operators $\hat{P}_N$ and $\hat{P}_{\pi}$: $|\Psi\rangle = \hat{P}_N |\Psi\rangle + \hat{P}_{\pi} |\Psi\rangle$. Then product of matrix elements in Eq.(13) in this notation is $\langle n_{\pi} | V | \hat{P}_N |\Psi\rangle$. Procedure of excluding of the meson channels gives rise to potential in $N\bar{N}$ -channel $\hat{P}_N U \hat{P}_{\pi} (H - E_{\pi})^{-1} \hat{P}_{\pi} V \hat{P}_N$ that is to compare with our short-range potential in Eq.(17). Here V denotes potential in the Shrödinger equation for the total wave function $|\Psi\rangle$.

5. Conclusion

We proposed the method of describing the $N\bar{N}$ -annihilation, based on the chiral soliton model. The crucial point was shown to be redefinition of the quantization procedure at small distances between the centres of a soliton and an antisoliton. This may be treated as the occurrence of the contribution to the nucleon – antinucleon potential. The radius of new interaction in our model is determined as a distance, where strong attraction violates the stability of the semiclassical quantization procedure in the soliton – antisoliton sector. Of course, this definition is rather rough but provides the usual in chiral models accuracy.

The next problem in evaluating the annihilation amplitude is the computation of the skyrmion – antiskyrmion wave function $\Psi_{N\bar{N}}(|\vec{R}| = R_0, \vec{\omega}, t_0)$. We have to solve the Shrödinger equation with classical potential $V_{N\bar{N}}$ and with boundary condition at the sphere $|\vec{R}| = R_0$. This condition arises from the fact that we have to "sew" the $N\bar{N}$ wave function with perturbative pion fields. It can be considered as resulting from the nonlocal potential (17) with additional restriction $\Psi_{N\bar{N}}(\vec{R})|_{|\vec{R}| < R_0} = 0$.

A large number of acts of the $N\bar{N}$ – decay into pions are known to go through formation of the resonance states. Our model can be easily extended to incorporate the description of vector mesons [17]. To include strangeness into the model one must consider $SU(3)$ - flavor group instead $SU(2)$ [18].

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References

- [1] M.P.Locher,B.S.Zou, Preprint PSI-PR-91-12, May 1991, Preprint PSI-PR-91-13, May 1991.
- [2] M.A.Alberg, E.M.Henley and L.Wilets, Z. Phys. **A331** (1988) 207.
- [3] S.Furui, Z. Phys. **C46** (1990) 621.
- [4] I.Zahed and G.E.Brown, Z. Phys. **A335** (1990) 349.
- [5] T.H.R.Skyrme, Nucl. Phys. **31** (1962) 556.
- [6] E.Witten, Nucl. Phys., **B223** (1983) 422, *ibid.* 433.
- [7] D.I.Dyakonov, V.Yu.Petrov, Nucl. Phys. **B272** (1986) 457.
- [8] C.Amsler, Preprint CERN/EP-86-178, Geneva, 1986.
- [9] B.O.Kerbikov, L.A.Kondrat'uk, M.G.Sapozhnikov, in: Proceedings of the XXIII "Particle Physics" LNPI Winter School, Leningrad, 1988, p.53.
- [10] J.Haidenbauer, T.Hippchen, K.Holinde, J.Speth, Z. Phys. **A334** (1989) 467.
- [11] J.-L.Gervais, A.Jevicki, Nucl. Phys. **B110** (1976) 93, *ibid.* 113.
- [12] M.M.Musakhanov, A.Rahimov, Preprint PTI-156-91-FVE, Tashkent, 1991.
- [13] I.I.Balitsky, A.V.Yung, Phys. Lett. **B168** (1986) 113; Nucl. Phys. **B274** (1986) 475.
- [14] A.Jevicki, Phys. Rev. **D21** (1980) 992.
- [15] R.P.Feynman, A.R.Hibbs. Quantum Mechanics and Path Integrals. NY: McGraw-Hill; 1965.
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- [16] A.M.Badalyan, M.I.Polikarpov, Yu.A.Simonov, Phys. Rep. **82** (1982) 31.
- [17] U.G.Meißner, I.Zahed, Phys. Rev. Lett. **56** (1986) 1035.
- [18] I.Zahed, G.E.Brown, Phys. Rep. **142** (1986) 1.

Figure captions

Figure 1:

The classical skyrmion - antiskyrmion potential V_{NN} [12] as a function of the relative distance R at various values of the orientation parameters:

a) $\vec{a}^2 = 0$, $T = 0$;

b) $\vec{a}^2 = 0.5$; solid line: $T^2 = 0$, dashed line: $T^2 = 0.5$, dotted line: $T^2 = 1.0$.

Fig. 1.

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