Abstract

The model of Kondo chain with $M$-fold degenerate band of conduction electrons of spin 1/2 interacting with localized spins $S$ is studied for the case when the electronic band is half filled. It is shown that the spectrum of spin excitations in the continuous limit is described by the O(3) nonlinear sigma model with the topological term with $\theta = \pi(2S - M)$. For a case $|M - 2S| = \text{(even)}$ the system is an insulator and single electron excitations at low energies are massive spin polarons. Otherwise the density of states has a pseudogap and vanishes only at the Fermi level. The relevance of this picture to higher dimensional Kondo insulators is discussed.

PACS numbers: 74.65.+n, 75.10. Jm, 75.25.+z

The problem of co-existence of delocalized and localized electrons in crystals remains one of the biggest unsolved problems in the condensed matter physics. The only part of this problem which is well understood concerns a situation when localized electrons are represented via a single local magnetic moment (the Kondo problem). In this case the local moment is screened at low temperatures by conduction electrons and the ground state is a singlet. The singlet formation is a non-perturbative process and the relevant energy scale (the Kondo temperature) is exponentially small in the exchange coupling constant. It is still unclear how conduction and localized electrons reconcile with each other when the local moments are arranged regularly (Kondo lattice problem). Empirically Kondo lattices resemble metals with very small Fermi energies of order of several degrees.
It is widely believed that conduction and localized electrons in Kondo lattices hybridize at low temperatures to create a single narrow band. It is not at all clear, however, how this hybridization develops. In particular, it is not clear whether the localized electrons contribute to the volume of Fermi sea. If the answer is positive, a system with one conduction electron and one spin per a unit cell must be an insulator. The available experimental data apparently support this point of view: all compounds with odd number of conduction electrons per spin are insulators\textsuperscript{1}. This class of compounds is now known as Kondo insulators. At low temperatures they behave as semiconductors with very small gaps of the order of several degrees. The marked exception is $FeSi$ where the value of gap is estimated as $\sim 700K$\textsuperscript{2}. The conservative approach to Kondo insulators would be to calculate their band structure treating the on-site Coulomb repulsion $U$ as a perturbation. The advantage of this approximation is that one gets an insulating state already in the zeroth order in $U$. The disadvantage is that it contradicts the principles of perturbation theory which prescribe that the strongest interactions are taken into account first. It turns out also that the pragmatic sacrifice of principles does not lead to a satisfactory description of the experimental data: the band theory fails to explain many experimental observations (see Ref. 2 for a discussion).

In this letter I study a one dimensional model of the Kondo lattice at the half filling. I show that the insulating state forms not due to a hybridization of conduction electrons with local moments, but as a result of strong antiferromagnetic fluctuations. Due to the local doubling of period of the original lattice the conduction electrons become heavier. An interaction of these heavy electrons with spin kinks lead to the formation of massive spin polarons. This scenario does not require a global antiferromagnetic order, just the contrary - the spin ground state remains disordered with a finite correlation length. I suggest that such scenario can be generalized for higher dimensions. Kondo insulators in this case are either antiferromagnets (then they have a true gap), or spin fluids with a strongly enhanced staggered susceptibility. In the latter case instead of a real gap there is a pseudogap - a drop in the density of states on the Fermi level. The recent numerical calculations of Yu et. al.\textsuperscript{3} also demonstrate a sharp enhancement of the staggered susceptibility in one-dimensional Kondo insulators.

As a model of one-dimensional Kondo insulator I consider the model of Kondo
chain at half filling governed by the following Hamiltonian:

\[ H = \sum_r \sum_{\alpha=1}^M \left[ -\frac{1}{2}(c_{r+1,\alpha,a}c_{r,\alpha,a} + c_{r,\alpha,a}c_{r+1,\alpha,a}) + J(c_{r,\alpha,a}^\dagger \vec{\sigma}_{ab}c_{r,\alpha,b})\right] \]  

(1)

It describes an \( M \)-fold degenerate band of electrons with spin \( S = 1/2 \) interacting with local spins \( S \). In what follows I shall use the path integral formalism. The path integral representation for spins has been discussed in details by many authors. I would refer a reader to the book of Fradkin\(^4\). In the path integral spins are treated as classical variables \( \vec{S} = S\vec{m} \) \((\vec{m}^2 = 1)\); the corresponding Euclidean action for the model (1) is given by:

\[
A = \int d\tau \left\{ \sum_r [iS \int_0^1 du (\vec{m}_r(u, \tau) \times \partial_\tau \vec{m}_r(u, \tau))] + c_{r,\alpha,a}^\dagger \partial_\tau c_{r,\alpha,a} - H(c^\dagger, c; S\vec{m}) \right\} 
\]

(2)

The first term is the spin Berry phase responsible for the correct quantization of local spins. Since the integrand in the Berry phase is a total derivative, the integral depends only on the value of \( \vec{m} \) on the boundary, i.e. on \( \vec{m}(u = 0, \tau) = \vec{m}(\tau), \vec{m}(u = 1, \tau) = (1, 0, 0) \). The introduction of the additional variable \( u \) is a price one has to pay for the fact that the Berry phase cannot be written as a local functional of \( \vec{m}_r(\tau) \).

I shall follow the semiclassical approach assuming that all fields can be separated into fast and slow components. The fast components then will be integrated out and as a result I shall obtain an effective action for the slow ones. This approach is self-consistent if the obtained correlation length for spins is much larger then the lattice spacing. In ordinary antiferromagnets this requirement is fulfilled only for large spins \( S >> 1 \). As we shall see later, the Kondo chain is semiclassical even for \( S = 1/2 \) provided the exchange integral is small \( JM << 1 \). I suggest the following decomposition of variables:

\[
\vec{m}_r = a\vec{k}(x) + (-1)^r\vec{n}(x)\sqrt{1 - a^2\vec{k}(x)^2}, \quad (\vec{k}\vec{n}) = 0, \\
c_r = i^r\psi_R(x) + (-i)^r\psi_L(x)
\]

(3)

where \(|\vec{k}|a << 1\) is the fastly varying ferromagnetic component of the local magnetization. Substituting Eqs.(3) into Eq.(2) and keeping only non-oscillatory
terms, I get:

\[ A = \int d\tau dx L, \]

\[ L = iS(\vec{k} [\vec{n} \times \partial_\tau \vec{n}]) + \bar{\psi}_j [i\gamma_\mu \partial_\mu \hat{I} + JS(\hat{\sigma} \vec{n}(x))\sqrt{1 - a^2 \vec{k}(x)^2}] \psi_j + 2\pi S \times (\text{top} - \text{term}) \] (4)

where

\[ \text{Top} - \text{term} = \frac{i}{8\pi} \int d\tau dx \epsilon_{\mu\nu} (\vec{n}[\partial_\mu \vec{n} \times \partial_\nu \vec{n}]) \] (5)

is the topological term first derived by Haldane\textsuperscript{5}. As far as the non-electronic part of the action is concerned, my derivation repeats the one presented in Ref. 4. The interaction of electrons with ferromagnetic fluctuations has been omitted; it can be shown that at small \( JM << 1 \) it gives insignificant corrections.

The fermionic determinant is calculated later (see Eq. (29)). Besides of the trivial static part it contains the topological term, but with \(-M\) instead of \(2S\). This is what one should expect: this change reflects the fact that local spins couple with conduction electrons to give the total spin \( S - M/2 \). Substituting the expression (30) into Eq. (4), I get:

\[ L = iS(\vec{k} [\vec{n} \times \partial_\tau \vec{n}]) + \frac{M}{2\pi} [(\partial_\tau \vec{n})^2 + (\partial_x \vec{n})^2] + \frac{2M}{\pi} (JS)^2 \ln \frac{1}{JS} \] (6)

Integrating over fast ferromagnetic fluctuations described by \( \vec{k} \), I get

\[ A = \frac{M}{2\pi} \int d\tau dx \left[ v^{-2} (\partial_\tau \vec{n})^2 + (\partial_x \vec{n})^2 \right] + \pi(2S - M) \times (\text{top} - \text{term}) \] (7)

\[ v^{-2} = 1 + \frac{2\pi^2}{J^2 M^2 \ln(1/JS)} \] (8)

After the rescaling of the coordinates \( v\tau = x_0, x = x_1 \) I get the action of the O(3) nonlinear sigma model with the dimensionless coupling constant

\[ g = \frac{\pi v}{M} = \frac{\pi}{\sqrt{M^2 + \frac{2\pi^2}{J^2 M^2 \ln(1/JS)}}} \] (9)

This constant is small at \( JM << 1 \) which justifies the entire semiclassical approach. At \(|M - 2S| = \) (even) one can omit the topological term. In this case the model Eq. (7) is the ordinary O(3) nonlinear sigma model. This model has a disordered ground state with the spectral gap\textsuperscript{6,7}

\[ \Delta = Jg^{-1} \exp[-2\pi/g] \] (10)
and the correlation length $\xi \sim J a / \Delta \gg a$. If $|M - 2S| = (\text{odd})$ the topological term is essential. The model becomes critical and the correlation functions of staggered magnetization have a power law decay. The specific heat is linear at small temperatures without requiring, however, the single electron density of states to be constant at the Fermi level.

Now I shall evaluate the fermionic determinant

$$D[g] = MT r \ln[i \gamma_\mu \partial_\mu + (1 + i\gamma_5) m g / 2 + (1 - i\gamma_5) m g^+ / 2]$$

where $g$ is a matrix from $SU(N)$ group and $m$ is some constant energy scale (in the context of the model $m = JS$ and $g = (\vec{\sigma} \vec{n})$). $N = 2$ in the original problem, but it is worth to do the calculation for general $N$. I shall study the expansion of the determinant in terms of $m^{-1} \nabla g$. The first terms of this expansion are independent of $m$ and survive even at $m \rightarrow \infty$. I claim that the gradient expansion contains a Berry phase. To prove this point I take a route which may seem exotic, but I do not know any better way to get the right answer. As a preliminary step I consider the chiral Gross-Neveu model with the $U(M) \times SU(N)$-symmetry described by the following action:

$$A = \int d^2 x \{ i \bar{\eta}_{a,\alpha} \gamma_\mu \partial_\mu \eta_{a,\alpha} - \frac{c}{2} [ (\bar{\eta}_{a,\alpha} \eta_{b,\alpha}) (\bar{\eta}_{b,\beta} \eta_{a,\beta}) - (\bar{\eta}_{a,\alpha} \gamma_5 \eta_{b,\alpha}) (\bar{\eta}_{b,\beta} \gamma_5 \eta_{a,\beta}) ] \}$$

The Greek indices belong to the group $SU(M)$ and the Latin ones to $SU(N)$. To avoid a confusion I emphasise that this model is not equivalent to the original model and I consider it only because the effective action for its low energy excitations are given by the determinant. In order to show that I introduce the auxiliary field $Q_{ab}$ and decouple the interaction term by the Hubbard-Stratonovich transformation. Interacting formally over the fermions we obtain the partition function for the tensor $Q_{ab}$:

$$Z = \int DQ^+ DQ \exp(- \int d^2 x L),$$

$$L = \frac{1}{2c} Tr Q^+ Q - M T r \ln[i \gamma_\mu \partial_\mu + (1 + i\gamma_5) Q / 2 + (1 - i\gamma_5) Q^+ / 2]$$

As in the standard $U(N)$-invariant Gross-Neveu model, the effective action has a saddle point with respect to $Q^+ Q$ and fluctuations of det $Q$ are massive. This can be shown in the standard fashion. Assuming that the saddle point configuration of $Q$ is coordinate independent and as such can be chosen as a diagonal real
matrix: $Q(x) = \text{diag}(\lambda_1, \ldots, \lambda_r)$, I calculate the density of effective action \((13)\) on this configuration:

$$A_{\text{eff}} = \sum_a \left[ \frac{\lambda_a^2}{2c} + \frac{M\lambda_a^2}{2\pi} \ln \left| \frac{\lambda_a}{\Lambda} \right| \right]$$

(14)

where $\Lambda$ is the ultraviolet cut-off. The saddle point value of $Q$, being a point of minimum of this function, satisfies the following equation:

$$\frac{\lambda_a}{c} + \frac{M\lambda_a}{\pi} \ln \left| \frac{\lambda_a}{\Lambda} \right| = 0$$

(15)

which solution is

$$\lambda_a = \Lambda \exp\left[-\frac{\pi}{Mc}\right] \equiv m$$

(16)

The performed calculation suggests that for slowly varying fields $Q$ one can substitute $Q_{ab}$ in the $\text{Tr} \ln$ in Eq. \((13)\) by $mg_{ab}$ where $g$ is an $SU(N)$ matrix. In other words, the effective action for excitations of the model \((12)\) with energies $<< m$ coincides with the fermionic determinant \((11)\) with $m$ given by Eq.\((14)\).

On the next step of the derivation I use the fact that due to the identity $2\tau_1^a \tau_2^a = 1/2 - P_{12}$, where $P_{12}$ is the permutation operator, the chiral Gross-Neveu model \((12)\) can be rewritten as the model with current-current interaction:

$$A = \int d^2 x \left\{ i\bar{\eta}_{a,\alpha} \gamma_\mu \partial_\mu \eta_{a,\alpha} + 4c J^\lambda_\mu J^\mu \lambda \right\}$$

(17)

where $\tau^r$ are matrices - generators of the $SU(M)$ group. The two models differ by a term containing a diagonal scattering. This term does not renormalize and therefore is not important. Now I apply to the model \((17)\) the non-Abelian bosonization procedure suggested by Witten\(^8\) (see also the book \(^9\)). Namely, I rewrite its Hamiltonian in the Sugawara form:

$$H = H_{U(1)} + H_{SU(N)} + H_{SU(M)},$$

(18)

$$H_{U(1)} = \pi \int dx \left[ J_R(x) J_R(x) : + : J_L(x) J_L(x) : \right]$$

(19)

$$H_{SU(N)} = \frac{2\pi}{(N+M)} \sum_{i=1}^{G_N} \int dx \left[ J^i_R(x) J^i_R(x) : + : J^i_L(x) J^i_L(x) : \right]$$

(20)

$$H_{SU(M)} = \sum_{\lambda=1}^{G_M} \int dx \left[ \frac{2\pi}{(N+M)} : J^\lambda_R(x) J^\lambda_R(x) : + : J^\lambda_L(x) J^\lambda_L(x) : + 4c J^\lambda_R(x) J^\lambda_L(x) : \right]$$

(21)

where I have introduced the chiral currents satisfying the Kac-Moody algebra; for the group $SU(N)$ the corresponding definition is

$$J^i_R = \eta_{R,aa}^i \bar{\eta}_{a,\alpha} \eta_{R,ba},$$

$$J^i_L = \eta_{L,aa}^i \bar{\eta}_{a,\alpha} \eta_{L,ba}.$$
\[ J^i_L = \eta^\dagger_{L,\alpha} t^i_{\alpha\beta} \eta_{L,\beta} \]  

where \( t^i \) are the generators of the SU(\( N \)) (spin) group and \( \eta_R, \eta_L \) are right and left components of the Dirac spinor \( \eta \). \( J \) and \( J^\lambda \) are the U(1) and the SU(M) (flavour) currents defined as in Eq. (22), but with generators of the corresponding algebras. \( G_N = N^2 - 1 \) and \( G_M = M^2 - 1 \) are the total number of generators of the \( su(N) \) and the \( su(M) \) algebras. Currents from different algebras commute.

Therefore the Hamiltonian (18) is a sum of three mutually commuting operators. Now notice that the interaction term containing right and left SU(M) currents do not affect the spectra of SU(M)-singlets. It is well known that the spectrum of \( H_{SU(M)} \) is gapful for the given sign of the coupling constant (see, for example, Ref. 10). From the previous discussion we can conclude that the gapful excitations correspond to fluctuations of \( \det Q \). Therefore the spectrum below the gap \( m \) is described by the rest of the Hamiltonian (18), in other words by

\[ H_{\text{eff}} = H_U(1) + H_{SU(N)}; \]  

\[ H_U(1) = \pi \int dx [J^i_R(x)J^i_R(x) : + : J^i_L(x)J^i_L(x) :]; \]  

\[ H_{SU(N)} = \int dx \frac{2\pi}{(N + M)} \sum_{a=1}^{G_N} (J^i_R(x)J^i_R(x) : + : J^i_L(x)J^i_L(x) :). \]  

The Hamiltonian (25) is the Hamiltonian of the Wess-Zumino-Witten model on the group SU(\( N \)). Its spectrum is the subsector of the free fermionic spectrum generated by the SU(\( N \)) current operators. The model is conformally invariant and exactly solvable\(^{11,12} \). In order to extract from these results an expression for the determinant (11), I rewrite the model (23) in the Lagrange representation\(^{8-12} \):

\[ A_{U(1)} = \frac{NM}{8\pi} \int d^2x (\partial_\mu \phi)^2 \]  

\[ A_{SU(N)} = \int d^2x \left\{ \frac{M}{16\pi} Tr(\partial_\mu g^+ \partial_\mu g) + \frac{M}{24\pi} \int_0^1 d\xi \epsilon_{abc} Tr(g^+ \partial_\alpha g g^+ \partial_\beta g g^+ \partial_\gamma g) \right\} \]  

where \( g \) is a matrix from the SU(\( N \)) group (\( g^+ g = I, \det g = 1 \)). The second term in the right hand side of Eq. (27) is called the Wess-Zumino term. This term is topological. Despite of the fact that it is written as an integral including the additional dimension, its actual value (modulus \( 2\pi iN \)) depends on the boundary values of \( g(x, \xi = 0) = g(x) \) (\( g(x, \xi = 1) = 0 \)). This property follows from the fact that the Wess-Zumino term is proportional to the integral of the Jacobian of transformation from the three dimensional euclidian coordinates to the group coordinates, and so it is a total derivative in disguise.
Now we can recognize in Eq. (27). Indeed, its first part (except of the Wess-Zumino term) represents the first term in the gradient expansion of $Tr \ln$-term in Eq. (13). Indeed, at small momenta $|p| << m$ this term is equal to

$$\frac{M}{16 \pi m^2} \int d^2 x Tr(\partial_\mu Q^+ \partial_\mu Q)$$

Now let us write down the $Q$-field as follows:

$$Q(x) = mg(x)e^{i\sqrt{M} \phi(x)},$$

where $g$ belongs to the $SU(N)$ group. Substituting this expression into Eq. (28) and taking into account that $Tr g^+ \partial_\mu g = 0$ as a trace of an element of the algebra, we reproduce Eq. (29) and the first term of Eq. (27). It comes not entirely unexpected that the naive gradient expansion misses the important Wess-Zumino term. Being a Berry phase this term requires special care and cannot be derived from a gradient expansion. In order to avoid possible calculational difficulties I have resorted to the exact solution which gives us the following interesting expression for the determinant (11):

$$D[g] = A_{SU(N)}[g] + \frac{M m^2}{2 \pi} \ln(\Lambda/m) + O(m^{-2})$$

where $A_{SU(N)}[g]$ is given by Eq. (27) and the second term represents the static part of the determinant. For the particular case $g = (\vec{\sigma} \vec{n})$ I get

$$A_{SU(2)}[\vec{\sigma} \vec{n}] = \frac{1}{2 \pi} (\partial_\mu \vec{n})^2 + \pi(\text{Top-term})$$

Let us discuss the excitation spectrum. I shall do it only for $M = 1$. In this case the original model has a combined symmetry $SU(2) \times SU(2)$ (the additional $SU(2)$-symmetry arises as a particle-hole symmetry at the half-filling), excitations carry two quantum numbers - spin $S$, and an isotopic spin $I$. We have established that spin excitations, i.e. excitations with $I = 0$, are described by the nonlinear sigma model with the topological term (4). The leading contributions to the low energy dynamics come from antiferromagnetic fluctuations which agrees with the results of Ref. 3. The corresponding energy scale (11) is formally resembles the expression for the Kondo temperature. $m(J)$ is larger, however, due to the presence of a large logarithm. Therefore the RKKY interaction plays a stronger role than the Kondo screening - it also agrees with the conclusions of Ref. 3. The topological term can be omitted if $|M - 2S| = (\text{even})$. In particular, it cancels for the most physical case $S = 1/2, M = 1$. The low lying magnetic excitations...
are in this case massive triplets, as it is for the O(3) nonlinear sigma model\textsuperscript{6,7}. This picture is in a qualitative agreement with the strong coupling limit of the model (\textsuperscript{[1]}). Indeed, for $J > > 1$ the ground state of the Kondo chain consists of local singlets. Excited states are triplets and are separated by the gap $\sim J$ from the ground state.

In order to describe the fermionic excitations which have $S = 1/2, I = 1/2$. It follows from (\textsuperscript{[4]}) that the fermionic fields live in a slowly fluctuating field $JS\{\vec{\sigma}\vec{n}(x, \tau)\}$. For constant $\vec{n}$ the electrons would have a spectral gap $JS$. There are states in the gap, however, and they are bound states of electrons and solitons of the unit vector field $\vec{n}$. Such situation is typical for relativistic field theories and for the $SU(N)$ chiral Gross-Neveu model it was discussed in Ref. 13. Let us consider a slow static configuration $\vec{n}(x)$ such that $\vec{n}(-\infty)$ is antiparallel to the $z$-axis and $\vec{n}(\infty)$ is parallel to it. Then a straightforward calculation shows that there is an electronic bound state on this domain wall with the zero energy. The corresponding wave function is equal to

$$\Psi(x) = \hat{T} \exp[-\int_{0}^{\infty} dy(\vec{\sigma}\vec{n}(y))]\Psi(0)$$

The total energy of this bound state is equal to the energy necessary to create the domain wall, which is the gap for the spin excitations. Therefore in the energy interval between $JS$ and $m$ single electron excitations are massive spin polarons. It is no longer the case if $|M - 2S| = (\text{odd})$. The nonlinear sigma model becomes critical and supposedly belongs to the universality class of the isotropic $S = 1/2$ Heisenberg chain. In the critical phase the spin solitons do not have a fixed scale and the corresponding bound states can have an arbitrary small energy. It is reasonable to suggest that the single particle density of states in this case has a pseudogap on the Fermi surface decaying as a power law: $\rho(\omega) \sim |\omega|^\alpha$.

We can generalize this one dimensional picture for higher dimensions. Suppose we have a Kondo lattice in three dimensions. Then at half filling we can achieve the insulating state by the periodicity doubling via an antiferromagnetic phase transition. Suppose, however, that this transition does not occur, but the system is very close to it. Then the low lying excitations are spin polarons. Whether they have a gap or not depends on the state of the spin system. For the one dimensional O(3) nonlinear sigma model the static magnetic susceptibility is strongly enhanced at the antiferromagnetic wave vector $q = \pi$ ($\chi(\pi) \sim m^{-1}$). In the same time it is zero at $q = 0$. The corresponding measurements in $Ce_3Bi_4Pt_3$ show a pronounced drop in $\chi(0)$ at low temperatures and the neutron measurements
show a spectral gap\(^1\). It indicates that the ground state in Kondo insulators is magnetically disordered. I postpone a discussion of three dimensional problem to further publications.

This work was started during my visit to the Department of Physics of Chalmers University in Goteborg (Sweden). I am grateful to Prof. S. Ostlund and the condensed matter group in Chalmers for the kind hospitality. I am grateful to A. Nersesyan and G. Japaridze for inspirational conversations, for P. Coleman for the valuable criticism, to Derek Lee for reading the manuscript and to C. Yu for sending me her preprint.

References

[1] G. Aeppli and Z. Fisk, Comments Condens. Matter Phys. 16, 155 (1992).

[2] Z. Schlesinger, Z. Fisk, Hai-Tao Zhang, M. B. Maple, J. F. DiTusa and G. Aeppli, Phys. Rev. Lett. 71, 1748 (1993).

[3] C. C. Yu, S. R. White, (unpublished).

[4] E. Fradkin, in "Field Theories of Condensed Matter Systems", Addison-Wesley Publ., Chapter 5.5, (1991).

[5] F. D. M. Haldane, Phys. Lett. 93A, 464 (1983); Phys. Rev. Lett. 50, 1153 (1983); J. Appl. Phys. 57, 3359 (1985).

[6] A. B. Zamolodchikov and Al. B. Zamolodchikov, Ann. Phys. (N.Y.) 120, 253 (1979).

[7] P. B. Wiegmann, Phys. Lett. 152B, 209 (1985).

[8] E. Witten, Comm. Math. Phys. 92, 455 (1984).

[9] C. Itzykson and J. - M. Drouffe, in Statistical Field Theory, v. 2, Appendix 9.C, Cambridge University Press Publ. (1989).

[10] A. M. Tsvelik, Zh. Eksp. Teor. Fiz. 93, 1329 (1987)[ Sov. Phys. JETP, 66, 754 (1987)].

[11] A. M. Polyakov and P. B. Wiegmann, Phys. Lett. B131, 121 (1983).

[12] V. G. Knizhnik and A. B. Zamolodchikov, Nucl. Phys. B247, 83 (1984).
[13] R. F. Dashen, B. Hasslacher and A. Neveu, Phys. Rev. D \textbf{12}, 2443 (1975).