Consistent regularization and renormalization in models with inhomogeneous phases

Prabal Adhikari\(^1\) and Jens O. Andersen\(^2,3\)

\(^1\)St. Olaf College, Physics Department, 1520 St. Olaf Avenue, Northfield, MN 55057, USA
\(^2\)Department of Physics, Faculty of Natural Sciences, NTNU, Norwegian University of Science and Technology, Høgskoleringen 5, N-7491 Trondheim, Norway
\(^3\)Niels Bohr International Academy, Blegdamsvej 17, Copenhagen 2100, Denmark

(Dated: August 3, 2017)

In many models in condensed matter and high-energy physics, one finds inhomogeneous phases at high density and low temperature. These phases are characterized by a spatially dependent condensate or order parameter. A proper calculation requires that one takes the vacuum fluctuations of the model into account. These fluctuations are ultraviolet divergent and must be regularized. We discuss different ways of consistently regularizing and renormalizing quantum fluctuations, focusing on momentum cutoff, symmetric energy cutoff, and dimensional regularization. We apply these techniques calculating the vacuum energy in the NJL model in 1+1 dimensions in the large-\(N_c\) limit and in the 3+1 dimensional quark-meson model in the mean-field approximation both for a one-dimensional chiral-density wave.

I. INTRODUCTION

There are many systems in condensed matter and high-energy physics where some of the phases are inhomogeneous. These are phases where an order parameter or a condensate depends on position. The simplest case is where only the phase of the order parameter is varying; the general case is where both magnitude and phase are functions of position. The idea of inhomogeneous phases is rather old going back to the work of Fulde and Ferrell as well as by Larkin and Ovchinikov in the context of superconductors \(^1\)\(^2\), density waves in nuclear matter by Overhauser \(^3\), and pion condensation by Migdal \(^4\). In recent years, inhomogeneous phases have been studied in for example cold atomic gases \(^5\), color superconducting phases \(^6\)–\(^8\), quarkyonic phases \(^9\)–\(^10\), as well as chiral condensates \(^11\)–\(^22\), see Refs. \(^23\)–\(^24\) for recent reviews.

In the case of QCD, these inhomogeneous phases exist at large baryon chemical potential \(\mu_B\) or isospin chemical potential \(\mu_I\), and low temperature. In the case of large \(\mu_B\), they can not be studied by lattice simulations due to the infamous sign problem and one must use low-energy models of QCD. Examples of such models are the Nambu-Jona-Lasinio (NJL) model and the quark-meson (QM) model or their Polyakov-loop extended versions, the PNJL and PQM models. Most of the calculations in 3+1 dimensions have been inspired by corresponding calculations in 1+1 dimensions using ansätze for the inhomogeneities that are one-dimensional \(^25\)–\(^33\). Interestingly, some of these models in 1+1 dimensions can be solved exactly in the large-\(N_c\) limit and show a rich phase diagram \(^28\)–\(^29\).

When calculating the thermodynamic potential in these models, one faces ultraviolet divergences due to vacuum fluctuations. The ultraviolet divergences in the NJL model are typically regularized using a sharp momentum cutoff \(\Lambda\) \(^34\). However, in the case of inhomogeneous condensates, a naïve application of a momentum cutoff leads to an incorrect expression for the vacuum energy in the limit where the magnitude of the order parameter vanishes \(^30\)–\(^31\). Instead, proper time regularization \(^12\), “symmetric energy cutoff regularization” \(^30\)–\(^31\) and Pauli-Villars regularization have been applied \(^13\)–\(^16\). In this paper, we will discuss how to use momentum cutoff, symmetric energy cutoff, and dimensional regularization in the case of inhomogeneous phases; in order to obtain a meaningful expression for the vacuum energy, it is necessary to perform a unitary transformation (that depends on the wave vector) on the free Hamiltonian and subtract the vacuum energy of the non-interacting system.

In the NJL model, one cannot throw away the quantum fluctuations since chiral symmetry breaking is induced by them, i.e. there is no symmetry breaking at tree level\(^1\). This is in contrast to the quark-
meson (QM) model, where the Higgs mechanism is implemented by chosing a negative mass term in the tree-level potential. In many finite-temperature applications of the quark-meson model, one ignores the vacuum fluctuations hoping that their effects on the chiral transition are negligible \[36\]. In the two-flavor QM model, the chiral transition is first order (in the chiral limit) in the whole $\mu_B-T$ plane without vacuum fluctuations. Adding the quantum fluctuations, the transition changes from being first order to being second order at zero baryon chemical potential $\mu_B$, while it remains first order at zero temperature. Thus, including the vacuum fluctuations, the first-order line that starts at $T=0$ ends at a tricritical point somewhere in the $\mu_B-T$ plane. Similarly, when allowing for an inhomogeneous phase such as a chiral-density wave, it exists in the entire $\mu_B-T$ plane in the absence of quantum fluctuations. Including quantum fluctuations, the inhomogeneous phase emerges from a tricritical point and exists in a region of low temperatures down to $T=0$ \[16\].

Conventionally, dimensional regularization has been used in the context of the quark-meson model. However, apriori, there is nothing that prevents us from treating the quark-meson model as a cutoff field theory \[33\] and regularizing it using a sharp ultraviolet momentum cutoff $\Lambda$. Having introduced an ultraviolet cutoff $\Lambda$, one can renormalize, i.e. redefine the parameters of the model and take the limit $\Lambda \rightarrow \infty$ at the end. In this case, one trades the ultraviolet cutoff for a renormalization scale $\mu$. In fact, this procedure yields results that are reminiscent of dimensional regularization in which power divergences are set to zero and logarithmic divergences show up as poles. The poles are then removed by renormalization of the parameters of the theory. Either way, there is an ambiguity, since there is a dependence on ultraviolet cutoff $\Lambda$ or the renormalization scale $\mu$.

The article is organized as follows. In Sec. II, we discuss the problem of a simple momentum cutoff in the context of an NJL type model in 1+1 dimensions. We will show that by subtracting the vacuum energy of the free theory after a unitary transformation, one can obtain a meaningful vacuum energy using a momentum cutoff, an energy cutoff, or dimensional regularization. In Sec. III, we show how to apply these techniques to the quark-meson model in three dimensions. In Sec. IV, we summarize and discuss our results.

II. NJL MODEL IN 1+1 DIMENSIONS

A. Lagrangian and thermodynamic potential

The Lagrangian of the NJL model in 1+1 dimensions is

\[
\mathcal{L} = \bar{\psi} \left( i \gamma^\mu \partial_\mu - m_0 + (\mu + \frac{1}{2} \tau_3 \mu_I) \gamma^0 \right) \psi \\
+ \frac{G}{N_c} \left( (\bar{\psi} \psi)^2 + (\bar{\psi} i \gamma^5 \tau_3 \psi)^2 \right),
\]

where $N_c$ is the number of colors and $m_0$ is the current quark mass. Moreover $\psi$ is a color $N_c$-plet, a two-component Dirac spinor as well as a flavor doublet

\[
\psi = \begin{pmatrix} u \\ d \end{pmatrix}.
\]

Here $\mu_B = 3\mu = \frac{1}{2}(\mu_u + \mu_d)$ and $\mu_I = (\mu_u - \mu_d)$ are the baryon and isospin chemical potentials expressed in terms of the quark chemical potentials $\mu_u$ and $\mu_d$. The $\gamma$-matrices are $\gamma^0 = \sigma_2$, $\gamma^1 = i\sigma_1$, and $\gamma^5 = \gamma^0 \gamma^1 = \sigma_3$, where $\sigma_i$ are the three Pauli matrices, and $\tau_a$ are the three Pauli matrices in flavor space. The Lagrangian \[1\] has a global $SU(N_c)$ symmetry and for $m_0 = \mu_I = 0$, it also invariant under $SU_L(2) \times SU_R(2)$. For $m_0 \neq 0$ and $\mu_I = 0$, the $SU_L(2) \times SU_R(2)$ symmetry is reduced to $SU_I(2)$. For $m_0 = 0$ and $\mu_I \neq 0$, the symmetry $SU_L(2) \times SU_R(2)$ is reduced to $U_{I_3 L}(1) \times U_{I_3 R}(1)$, where $I_3$ is the third component of isospin. If $m_0 \neq 0$ and $\mu_I \neq 0$ the $SU_L(2) \times SU_R(2)$ symmetry is reduced to $U_{I_3}(1)$.

We next introduce the collective sigma and pion fields

\[
\sigma = -\frac{2G}{N_c} \bar{\psi} \psi, \\
\pi_a = -\frac{2G}{N_c} \bar{\psi} i \gamma^5 \tau_a \psi.
\]

The Lagrangian \[1\] then becomes

\[
\mathcal{L} = \bar{\psi} \left( i \gamma^\mu \partial_\mu - m_0 + (\mu + \frac{1}{2} \tau_3 \mu_I) \gamma^0 \right) \psi \\
- \frac{N_c (\sigma^2 + \pi_a^2)}{4G}.
\]

The chiral condensate that we choose is a chiral-density wave of the form\[2\]

\[
\langle \sigma \rangle = M \cos(2bz) - m_0, \\
\langle \pi_3 \rangle = M \sin(2bz),
\]

2 With a nonzero isospin chemical potential, there is also the possibility of a pion condensate $\Delta$. For simplicity, we do not include this in the present analysis.
where \( b \) is a wavevector.

The Mermin-Wagner-colemann theorem normally forbids spontaneous symmetry breaking in 1+1 dimensions; however it does not apply in the large \( N_c \) limit \[38, 39\]. We denote the last term in Eq. (5) by \( V_0 \) such that \( V_0 \) is the tree-level potential. Note that for nonzero \( z \), the crossterm \(-N_c m \cos(2bz)/2G\) averages to zero when the spatial extent \( L \) of the system large enough. This term can then be written as \[-N_c m \cos(2b_{k=0})/2G\] and the expression for \( V_0 \) is

\[
V_0 = \frac{N_c (M^2 + m^2 - 2M m_0 \delta_{k=0})}{4G}.
\]

In the homogeneous case, \( V_0 = \frac{N_c (M - m_0)^2}{4G} \).

With the ansatz (6)–(7), the Dirac operator \( D \) can be written as

\[
D = \bar{\psi} \left[ i \frac{\partial}{\partial \tau_3} + (\mu + b\tau_3 \mu_I) \gamma^0 - M e^{2i\gamma^5 \tau_3 b_z} \right] \psi.
\]

We next redefine the quark fields, \( \psi \rightarrow e^{-i\gamma^5 \tau_3 b_z} \psi \) and \( \bar{\psi} \rightarrow \bar{\psi} e^{i\gamma^5 \tau_3 b_z} \). The Dirac operator then reads

\[
D = \bar{\psi} \left[ i \frac{\partial}{\partial \tau_3} + (\mu + b' \tau_3) \gamma^0 - M \right],
\]

where \( b' = (b + \tau_3 \mu_I) \) and \( 2b' \) is an effective isospin chemical potential. The transformation of the field \( \psi \) amounts to a unitary transformation of the Dirac Hamiltonian, \( \mathcal{H} \rightarrow \mathcal{H}' = e^{i\gamma^5 \tau_3 b_z} \mathcal{H} e^{-i\gamma^5 \tau_3 b_z} \). It turns out that there is a spurious dependence on \( b \) in the free energy: For some regulators, the free energy depends on \( \mu_I \). However, physical quantities cannot depend on the wavevector when the modulus of the condensate is zero. This unphysical behavior of the free energy requires the introduction of a subtraction term that by construction guarantees that the free energy is independent of \( b \) in the limit \( M \rightarrow 0 \). We will return to this issue below. There is an additional complication for nonzero isospin chemical potential since the spurious dependence of \( b \) translates into additional dependence on the isospin chemical potential in the free energy. We therefore set \( \mu_I = 0 \) for now and return to the case of nonzero \( \mu_I \) at the end of this section.

Going to momentum space, Eq. (10) can be written as

\[
D = \left[ p + (\mu + b \tau_3) \gamma^0 - M \right].
\]

It is now straightforward to derive the fermionic spectrum in the background (7). It is given by the zeros of the Dirac determinant and reads \([30]\)

\[
E_{\pm} = \sqrt{(p^2 + M^2 \pm b)^2}.
\]

We notice that the lower branch, \( E_- \), has zero energy for nonzero momentum, \( p = \pm \sqrt{b^2 - M^2} \), if \( b > M \). It is this nonmonotonic behavior that allows for inhomogeneous condensates at finite chemical potential by lowering the energy and at the same time populating only the lower branch \( E_- \).

We can now integrate over the fermions to obtain the free energy in the mean-field approximation. After integrating over \( p_0 \), this yields the standard expression

\[
V = V_0 - N_c \int_{-\infty}^{\infty} \frac{dp}{2\pi} (E_+ + E_-). \tag{13}
\]

### B. Momentum versus energy cutoff

The starting point is the one-loop correction to the effective potential,

\[
V_1 = -N_c \int_{-\infty}^{\infty} \frac{dp}{2\pi} (E_+ + E_-). \tag{14}
\]

If we regulate the integral by a simple momentum cutoff \( \Lambda \), we can write

\[
V_1 = -\frac{N_c}{\pi} \int_{0}^{\Lambda} (E_+ + E_-) dp, \tag{15}
\]

It will prove useful to change variable to \( u = \sqrt{p^2 + M^2} \). The integral then becomes

\[
V_1 = -\frac{N_c}{\pi} \int_{M}^{\Lambda} |u + b| + |u - b| \frac{udu}{\sqrt{u^2 - M^2}}. \tag{16}
\]

We next write \( V_1 = V_+ + V_- \). In the limit of large \( \Lambda \), we find

\[
V_+ = -\frac{N_c}{\pi} \int_{M}^{\Lambda} \sqrt{u^2 + M^2} |u + b| \frac{udu}{\sqrt{u^2 - M^2}} = -\frac{N_c}{2\pi} \left[ \Lambda^2 + 2b \Lambda + \frac{1}{2} M^2 \left( \log \frac{4\Lambda^2}{M^2} + 1 \right) \right]. \tag{17}
\]

The other contribution \( V_- \) is given by

\[
V_- = -\frac{N_c}{\pi} \int_{M}^{\Lambda} \sqrt{u^2 + M^2} |u - b| \frac{udu}{\sqrt{u^2 - M^2}}. \tag{18}
\]

Here we must be careful distinguishing between the cases \( u > b \) and \( u < b \). In the large-\( \Lambda \) limit, one finds

\[
V_- = -\frac{N_c}{2\pi} \left[ \Lambda^2 - 2b \Lambda + \frac{1}{2} M^2 \left( \log \frac{4\Lambda^2}{M^2} + 1 \right) \right] + \theta(b - M) f(M,b), \tag{19}
\]

where we have defined the function \( f(M,b) \).
The one-loop contribution to the free energy is then given by the sum of Eqs. (17) and (19). After renormalizing the vacuum energy by removing the term proportional to $\Lambda^2$, the effective potential in the mean-field approximation becomes

$$V = V_0 - \frac{N_c M^2}{2\pi} \left[ \log \frac{4\Lambda^2}{M^2} + 1 \right] + \theta(b - M) f(M, b) .$$

We note that the terms linear in $b$ cancel and that the final result is an even function of $b$ as it must, cf. Eq. (16). However, the vacuum energy is unbounded below due to the term $b\sqrt{b^2 - M^2}$ implying that the system is unstable. Moreover, in the limit $M \to 0$, the effective potential reduces to $V = -\frac{N_c b^2}{2\pi}$ (for $m_0 = 0$). This is clearly unphysical; the effective potential must be independent of the wavevector $b$ when the magnitude $M$ of the condensate vanishes. As pointed out in [30, 31], the problem is that the cutoff is imposed on the momentum and not the energy. Using a momentum cutoff $\Lambda$, the effective cutoff on the energy is $\sqrt{\Lambda^2 + M^2} \pm b$, which is different for the two branches of nonzero $b$. The idea put forward in [30, 31] is to use a cutoff $\Lambda$ on the energy rather than the momentum of the particles, i.e. one restricts the integration by imposing the same cutoff on the two branches,

$$E_{\pm} < \Lambda .$$

This is referred to as symmetric energy cutoff [30, 31]. This restriction can be expressed as an upper limit for the integration variable $u$ and yields $u < \Lambda \mp b$. The contribution $V_+$ from $E_+$ is in the large-$\Lambda$ limit then becomes

$$V_+ = - \frac{N_c}{\pi} \int_M^{\Lambda - b} |u + b| \frac{udu}{\sqrt{u^2 - M^2}}$$

$$= - \frac{N_c}{2\pi} \left[ \Lambda^2 - b^2 + \frac{1}{2} M^2 \left( \log \frac{4\Lambda^2}{M^2} - 1 \right) \right] .$$

The contribution $V_-$ from $E_-$ in the large-$\Lambda$ limit given by

$$V_- = - \frac{N_c}{\pi} \int_M^{\Lambda + b} |u - b| \frac{udu}{\sqrt{u^2 - M^2}}$$

$$= - \frac{N_c}{2\pi} \left[ \Lambda^2 - b^2 + \frac{1}{2} M^2 \left( \log \frac{4\Lambda^2}{M^2} - 1 \right) \right]$$

$$+ \theta(b - M) f(M, b) .$$

Note that $V_+ = V_-$ when $M > b$. The final result for the vacuum energy is given by the sum of the tree-level term, Eqs. (23) and (24)

$$V = V_0 - \frac{N_c}{2\pi} \left[ -2b^2 + M^2 \left( \log \frac{4\Lambda^2}{M^2} - 1 \right) \right]$$

$$+ \theta(b - M) f(M, b) ,$$

where we again have added the vacuum counterterm $\frac{N_c \Lambda^2}{2\pi}$. In the limit $M \to 0$, the $b$-dependent terms in Eq. (25) drop out and the effective potential vanishes (again for $m_0 = 0$). Thus the symmetric-energy cutoff provides us with a well-defined effective potential.

Returning to the momentum cutoff, one can of course simply subtract the term $-\frac{N_c b^2}{2\pi}$ from the effective potential, but this requires some justification. The idea is to subtract the vacuum energy of the noninteracting system, i.e. that of a free Fermi gas, as a part of the renormalization prescription [16]. As explained above, the redefinition of the quark fields immediately after Eq. (9) corresponds to a unitary transformation of the Hamiltonian of the system. We must therefore also perform the same transformation on the free Hamiltonian. The subtraction term is then obtained by making the substitution $M \to m_0$ in Eq. (21). The total effective potential is given by the sum of

$$V = V_0 - \frac{N_c M^2}{2\pi} \left[ \log \frac{4\Lambda^2}{M^2} + 1 \right] + \theta(b - M) f(M, b)$$

$$- \theta(b - m_0) f(m_0, b) ,$$

where we have dropped terms that depend on $\Lambda$ and $m_0$. Taking the limits $m_0 \to 0$ and $M \to 0$ (in this order), we see that all the $b$-dependent terms cancel as they should. We note in passing that the subtraction term is not unique. We could have subtracted the vacuum energy of a massless Fermi gas, $V_{\text{sub}} = -\frac{N_c b^2}{2\pi}$, which was done in Ref. [30]. If we use this prescription, the expression for the vacuum energy becomes

$$V = V_0 - \frac{N_c}{2\pi} \left[ -2b^2 + M^2 \left( \log \frac{4\Lambda^2}{M^2} + 1 \right) \right]$$

$$+ \theta(b - M) f(M, b) .$$
C. Dimensional regularization

Let us next consider the vacuum energy using dimensional regularization. The one-loop energy is given by the sum of the two terms

\[ V_\pm = -N_c \int_p E_\pm , \tag{28} \]

where the integral is defined in \( d = 1 - 2\epsilon \) dimensions:

\[ \int_p = \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right) \epsilon \int \frac{d^d p}{(2\pi)^d}. \tag{29} \]

Here \( \Lambda \) is the renormalization scale associated with the \( \overline{\text{MS}} \) renormalization scheme. We first change variables \( u = \sqrt{p^2 + M^2} \) and integrate over angles,

\[
V_\pm = -N_c \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right) \epsilon \int M \left( \frac{|u-b| u du}{(u^2 - M^2)^{1/2 + \epsilon}} \right).
\]

This yields

\[
V_\pm = -N_c \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right) \epsilon \int M \left( \frac{|u-b| u du}{(u^2 - M^2)^{1/2 + \epsilon}} \right).
\]

We first consider the contribution \( V_\pm \) to the effective potential from \( E_\pm \). We find

\[
V_\pm = \frac{N_c M^2}{4\pi} \left( \frac{e^{\gamma_E} \Lambda^2}{M^2} \right) \Gamma(-1 + \epsilon). \tag{31} \]

The contribution \( V_\pm \) is independent of \( b \), which is most easily understood by going back to momentum space in Eq. (30). The \( b \)-dependence is then given by an integral over \( p \) with no mass scale multiplied by \( b \), and this integral vanishes in dimensional regularization.

The contribution \( V_- \) from the negative solution \( E_- \) is given by

\[
V_- = -N_c \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right) \epsilon \int M \left( \frac{|u-b| u du}{(u^2 - M^2)^{1/2 + \epsilon}} \right) + 2\theta(b - M) \int M \left( \frac{(b-u) u du}{(u^2 - M^2)^{1/2 + \epsilon}} \right).
\]

where the second integral has been evaluated directly in one dimension. We note that the first term in Eq. (32) is equal to \( V_+ \).

The total vacuum energy is given by the sum of the tree-level term, Eqs. (31) and (32). Expanding this expression in powers of \( \epsilon \), we obtain

\[
V = V_0 - \frac{N_c M^2}{2\pi} \left( \frac{\Lambda^2}{M^2} \right) \epsilon \left[ \frac{1}{\epsilon} + 1 \right] + \theta(b - M) f(M, b). \tag{33} \]

The pole in \( \epsilon \) is removed by renormalizing the quark mass \( m_0 \) and the constant coupling \( G \). This is carried out by the substitutions \( m_0 \to Z_{m_0} m_0 \), and \( \frac{1}{G} \to Z_{G^{-1}} \frac{1}{G} \) where the mass and inverse coupling renormalization constants are

\[
Z_{m_0} = \left[ 1 + \frac{2G}{\pi\epsilon} \right]^{-1}, \quad \tag{34} \]

\[
Z_{G^{-1}} = \left[ 1 + \frac{2G}{\pi\epsilon} \right]. \quad \tag{35} \]

Note that \( Z_{G^{-1}} = Z_G^{-1} \) and that the ratio \( \frac{m_0}{m} \) is the same for bare and renormalized quantities since

\[
Z_{m_0} Z_{G^{-1}} = 1. \]

This yields the renormalized effective potential in the mean-field approximation

\[
V = V_0 - \frac{N_c M^2}{2\pi} \left[ \log \left( \frac{\Lambda^2}{M^2} \right) + 1 \right] + \theta(b - M) f(M, b). \tag{36} \]

We note in passing that the substitutions (34)–(35) correspond to a nonperturbative renormalization. In perturbation theory, it amounts to summing an infinite series of diagrams from all orders of perturbation theory. This can for example be seen by analyzing the model in terms of the two-particle irreducible effective action formalism to leading order in the \( 1/N_c \)-expansion in analogy with the bosonic case in three dimensions. Moreover, the running coupling \( G \) and the running mass \( m_0 \) satisfy the renormaliza-
tion group equations
\[ \Lambda \frac{dG}{d\Lambda} = -\frac{4G^2}{\pi}, \quad (37) \]
\[ \Lambda \frac{d\mu_0}{d\Lambda} = -\frac{4\mu_0 G}{\pi}. \quad (38) \]

However, Eq. (38) is still problematic. Taking the limit \( M \to 0 \), we find \( V = -\frac{N_c b^2}{2} \) (for \( m_0 = 0 \)) which is unphysical. In order to understand the source of the problem, we must go back to the contribution \( V_\pm \) and take the limit \( M \to 0 \):
\[ V_\pm = -\frac{N_c}{\pi} \left( \frac{e^{b^2\Lambda^2}}{4\pi} \right)^\epsilon \int_0^\infty |p \pm b| p^{-2\epsilon} dp. \quad (39) \]

\( V_\pm \) can be written as a sum of two integrals in which there is no mass scale. These integrals are then set to zero in dimensional regularization. In \( V_\pm \), the wavevector \( b \) is the only scale in the integral and according to the rules of dimensional regularization, the integral will be proportional to the appropriate power of \( b \). Dimensional analysis gives \( V_\pm \sim b^{2-2\epsilon} \). The coefficient is finite and in the limit \( \epsilon \to 0 \) one finds \( V_\pm = -\frac{N_c b^2}{2} \). This expression is exactly the vacuum energy of \( N_c \) massless fermions after a unitary transformation of the Hamiltonian. As in case of a momentum cutoff, we subtract the vacuum energy in the normal phase as part of renormalization procedure. This is given by the second and third term in Eq. (38) after the substitution \( M \to m_0 \). Dropping trivial terms that depend on \( m_0 \) and \( \Lambda \), we find
\[ V = V_0 - \frac{N_c M^2}{2\pi} \left[ \log \frac{\Lambda^2}{M^2} + 1 \right] + \theta(b - M) f(M, b) - \theta(b - m_0) f(m_0, b). \quad (40) \]

We close this section by noting that in the chiral limit, the gap equation \( \frac{dV}{dM} = 0 \), in the vacuum where \( b = 0 \), has two solutions, either \( M_0 = 0 \) or
\[ M_0 = \Lambda e^{-\frac{a}{4\pi}}. \quad (41) \]

Using the renormalization group equation (38), it is easy to verify that \( M_0 \) is renormalization group invariant. The nonanalytic behavior of \( M_0 \) as a function of \( G \) shows that it is a nonperturbative result. As mentioned above, it corresponds to summing an infinite series of diagrams from all order of perturbation theory.

Eq. (41) can be used to trade the cutoff or the renormalization scale for the mass scale \( M_0 \). In dimensional regularization, we find
\[ V = -\frac{N_c M^2}{2\pi} \left[ \log \frac{M_0^2}{M^2} + 1 \right]. \quad (42) \]

The RG-invariance of \( M_0 \) implies the RG-invariance of \( V \) in (42). As pointed out in 30, the unrenormalized expression for the vacuum energy contains a dimensionless parameter \( G \), while the renormalized result (42) contains a dimensionful mass scale \( M_0 \). This is an example of dimensional transmutation.

III. QUARK-MESON MODEL

A. Lagrangian and thermodynamic potential

The Euclidean Lagrangian of the two-flavor quark-meson model is
\[ \mathcal{L} = \frac{1}{2} \left[ (\partial_\mu \sigma)^2 + (\partial_\mu \pi)^2 \right] + \frac{1}{2} m^2 (\sigma^2 + \pi^2) - h \sigma + \frac{\Lambda}{24} (\sigma^2 + \pi^2)^2 + \bar{\psi} \left[ i \gamma^\mu (\mu - \frac{1}{2} \tau_3 \mu_1) \gamma^\mu + g(\sigma + i\gamma^5 \tau \cdot \pi) \right] \psi, \]
\[ \psi = \begin{pmatrix} u \\ d \end{pmatrix}. \quad (43) \]

where \( \psi \) is a color \( N_c \)-plet, a four-component Dirac spinor as well as a flavor doublet

\[ \psi = \begin{pmatrix} u \\ d \end{pmatrix}. \quad (44) \]

Here \( \mu_B = 3 \mu = \frac{3}{2} (\mu_u + \mu_d) \) and \( \mu_I = (\mu_u - \mu_d) \) are the baryon and isospin chemical potentials expressed in terms of the quark chemical potentials \( \mu_u \) and \( \mu_d \), \( \tau_i \) (\( i = 1, 2, 3 \)) are the Pauli matrices in flavor space, and \( \tau = (\pi_1, \pi_2, \pi_3) \).

Apart from the global \( SU(N_c) \) symmetry, the Lagrangian (13) has a \( U(1)_B \times SU(2)_L \times SU(2)_R \) symmetry for \( h = 0 \) and a \( U(1)_B \times SU(2)_V \) symmetry.
for \( h \neq 0 \). When \( \mu_u \neq \mu_d \), this symmetry is reduced to \( U(1)_B \times U(1)_L \times U(1)_R \) for \( h = 0 \) and \( U(1)_B \times U(1)_L \) for \( h \neq 0 \). In the remainder of this paper we take \( h = 0 \), i.e. we work in the chiral limit. We also set \( \mu_u = \mu_d \).

In order to study inhomogeneous phases, we must make an ansatz for the space-time dependence of the mesonic mean fields. In the literature, mainly one-dimensional modulations have been considered, for example chiral-density waves (CDW) and soliton lattices. Since the results seem fairly independent of the modulation \( [16] \), we opt for the simplest, namely a chiral-density wave. The ansatz is

\[
\sigma(z) = \phi_0 \cos(qz), \quad \pi_3(z) = \phi_0 \sin(qz),
\]

where \( \phi_0 \) is the magnitude of the condensate and \( q \) is the wavevector. The mean fields can be be combined into a complex order parameter as

\[
M(z) = g[\sigma(z) + i\pi_3(z)]
\]

or

\[
M(z) = \Delta e^{iqz},
\]

where we have introduced \( \Delta = g\phi_0 \). After averaging over a sufficiently large volume \( V_3 \) in three dimensions, the tree-level effective potential is then

\[
V_0 = \frac{1}{2}q^2 \frac{\Delta^2}{g^2} + \frac{1}{2}m^2 \frac{\Delta^2}{g^2} + \frac{\lambda}{24} \frac{\Delta^4}{g^4}.
\]

In analogy with the previous example, we can derive the spectrum by finding the zeros of the Dirac determinant. The result is \( [43] \)

\[
E_{\pm} = \sqrt{\left( \sqrt{p^2_{\|} + \Delta^2} \pm \frac{q}{2} \right)^2 + p^2_{\perp}},
\]

where \( p^2_{\|} = p^2_1 + p^2_2 \), and \( p_{\|} = p_1 \). Note that the lower branch has a vanishing minimum, \( E_- = 0 \), for nonzero momentum \( p_{\|} = \pm \sqrt{\frac{q^2}{4} - \Delta^2} \) and \( p_{\perp} = 0 \) in the case \( \frac{q}{2} > \Delta \). It is this nonmonotonic behavior that allows for inhomogeneous condensates at finite density; it may be energetically favorable for the system to develop a nonzero value of \( q \) and populate only the lower branch \( E_- \). Although inhomogeneous phases are possible only for nonzero chemical potentials, the vacuum energy is independent of \( \mu_f \) so the chemical potentials play no role in the calculation below.

### B. Energy and momentum cutoff

The vacuum part of the one-loop contribution to the effective potential is given by the expression

\[
V_1 = -2N_c \int (E_+ + E_-),
\]

where the integral is in three spatial dimensions. We first use an energy cutoff to evaluate Eq. \( [49] \). In the case of \( E_+ \), we must distinguish between the cases \( \sqrt{p^2_{\|} + \Delta^2} - \frac{q}{2} > 0 \), and \( \sqrt{p^2_{\|} + \Delta^2} - \frac{q}{2} < 0 \). As in 1+1 dimensions, there is an extra term in the case \( \sqrt{p^2_{\|} + \Delta^2} - \frac{q}{2} < 0 \), which we denote by \( f(\Delta, q) \). We first integrate over \( p_{\perp} \) from zero to \( p_{\perp}^{\max} = \sqrt{\Delta^2 - \left( \frac{u \pm \frac{q}{2} \right)^2} \), and then integrate over \( u = \sqrt{p^2_{\|} + \Delta^2} \) from \( u = \Delta \) to \( u = \Lambda \mp \frac{q}{2} \) (upper sign for \( E_+ \) and lower sign for \( E_- \)). The expressions for the integrals are
\[ V_+ = -2N_c \int_p E_+ \]
\[ = - \frac{2N_c}{(4\pi)^2} \left[ \sqrt{\frac{(\Lambda - \frac{q}{2})^2 - \Delta^2}{\Delta}} \left( \left( \frac{q}{2} \right)^2 - \Delta^2(6\Lambda + 13q) \right) \right] \]
\[ - (\Delta^4 + \Delta^2q^2) \log \left( \frac{\Lambda - \frac{q}{2} + \sqrt{(\Lambda - \frac{q}{2})^2 - \Delta^2}}{\Delta} \right) \right] , \tag{50} \]
\[ V_- = -2N_c \int_p E_- \]
\[ = - \frac{2N_c}{(4\pi)^2} \left[ \sqrt{\frac{(\Lambda + \frac{q}{2})^2 - \Delta^2}{\Delta}} \left( \left( \frac{q}{2} \right)^2 - \Delta^2(6\Lambda - 13q) \right) \right] \]
\[ - (\Delta^4 + \Delta^2q^2) \log \left( \frac{\Lambda + \frac{q}{2} + \sqrt{(\Lambda + \frac{q}{2})^2 - \Delta^2}}{\Delta} \right) \right] + \theta(\frac{q}{2} - \Delta)f(\Delta, q) , \tag{51} \]

where the function \( f(\Delta, q) \) is defined as
\[ f(\Delta, q) = \frac{N_c}{3(4\pi)^2} \left[ q\sqrt{\frac{q^2}{4} - \Delta^2(26\Delta^2 + q^2) - 12\Delta^2(\Delta^2 + q^2) \log \frac{q + 2\sqrt{\frac{q^2}{4} - \Delta^2}}{2\Delta} \right] . \tag{52} \]

In the limit \( \Delta \to 0 \), \( V_+ + V_- \) reduces to \( -\frac{8N_c}{(4\pi)^2}A^4 \) showing that the thermodynamic potential is independent of \( q \) in this limit. Subtracting this term then corresponds to a trivial renormalization of the vacuum energy. In the limit \( \Lambda \to \infty \), the sum of \( \text{(50)} \) and \( \text{(51)} \) behaves as
\[ V_+ + V_- = - \frac{2N_c}{(4\pi)^2} \left[ -4\Lambda^2\Delta^2 - q^2\Delta^2 \left( \log \frac{4\Lambda^2}{\Delta^2} - 2 \right) - \Delta^4 \left( \log \frac{4\Lambda^2}{\Delta^2} - \frac{1}{2} \right) + \frac{1}{12}q^4 \right] + f(\Delta, q) , \tag{53} \]
in agreement with the result first obtained by Broniowski and Kutschera \[43\].

Let us briefly discuss the calculation of the vacuum energy using a momentum cutoff \( \Lambda \). Integrating Eqs. \( \text{(50)} \) and \( \text{(51)} \) and taking the limit \( \Delta \to 0 \), we find \( V_+ + V_- = -\frac{8N_c}{(4\pi)^2}(\Lambda^4 + \frac{2}{3}q^2\Lambda^2 + \frac{\pi^2}{3}q^4) \), which must be subtracted. For large \( \Lambda \), the final result is
\[ V_+ + V_- = - \frac{2N_c}{(4\pi)^2} \left[ 4\Lambda^2\Delta^2 - q^2\Delta^2 \left( \log \frac{4\Lambda^2}{\Delta^2} - \frac{5}{3} \right) - \Delta^4 \left( \log \frac{4\Lambda^2}{\Delta^2} - \frac{1}{2} \right) + \frac{1}{12}q^4 \right] + f(\Delta, q) . \tag{54} \]

Comparing Eqs. \( \text{(53)} \) and \( \text{(54)} \), we see that the coefficients of some of the terms are different. However, the coefficients of the logarithmic terms are identical.

\[ \int_p = \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right)^\epsilon \int \frac{d^d p}{(2\pi)^d} = \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right)^\epsilon \int_{p_\perp} \frac{d^{d-1} p_\perp}{(2\pi)^{d-1}} \int_{p_\parallel} \frac{d p_\parallel}{2\pi} . \tag{56} \]

\[ V_\pm = -2N_c \int_p E_\pm , \tag{55} \]

We first integrate over angles in the \( (p_1, p_2) \)-plane and introduce the variable \( u = \sqrt{p_1^2 + \Delta^2} \). The inte-
We denote the right-hand side of (58) by \( \text{sub} \). We next integrate \( V \). This yields

\[
\sqrt{(u \pm \frac{q}{2})^2 + p_\perp^2} = \sqrt{u^2 + p_\perp^2} + \frac{uq}{2\sqrt{u^2 + p_\perp^2}} + \frac{q^2 p_\perp^2}{8(u^2 + p_\perp^2)^{\frac{3}{2}}} - \frac{q^4 p_\perp^2}{16(u^2 + p_\perp^2)^{\frac{5}{2}}} + \frac{q^4 p_\perp^2 (4u^2 - p_\perp^2)}{128(u^2 + p_\perp^2)^{\frac{7}{2}}} + \ldots (58)
\]

We denote the right-hand side of (58) by \( \text{sub} (u, p_\perp) \) and write the integrals in (57) as

\[
V_\pm = V_{\text{div}, \pm} + V_{\text{fin}, \pm} - V_{\text{fin}, \pm},
\]

where

\[
V_{\text{div}, \pm} = -\frac{N_c (e^\gamma A^2)^\epsilon}{\pi^2 \Gamma(1 - \epsilon)} \int_\Delta^\infty \frac{u \, du}{\sqrt{u^2 - \Delta^2}} \int_0^\infty \text{sub}_\pm (u, p_\perp) p_\perp^{1-2\epsilon} dp_\perp,
\]

\[
V_{\text{fin}, \pm} = -\frac{N_c (e^\gamma A^2)^\epsilon}{\pi^2 \Gamma(1 - \epsilon)} \int_\Delta^\infty \frac{u \, du}{\sqrt{u^2 - \Delta^2}} \int_0^\infty \left[ \sqrt{(u \pm \frac{q}{2})^2 + p_\perp^2} - \text{sub}_\pm (u, p_\perp) \right] p_\perp^{1-2\epsilon} dp_\perp.
\]

The integral \( V_{\text{fin}, \pm} \) can now be calculated directly in three dimensions. After integrating over \( p_\perp \), we find

\[
V_{\text{fin}, \pm} = -\frac{N_c}{3\pi^2} \int_\Delta^\infty \frac{u \, du}{\sqrt{u^2 - \Delta^2}} (u \pm \frac{q}{2})^2 \left( (u \pm \frac{q}{2})^2 - |u \pm \frac{q}{2}| \right).
\]

Thus \( V_{\text{fin}, -} \) vanishes identically and \( V_{\text{fin}, +} \) becomes

\[
V_{\text{fin}, +} = -\frac{2N_c}{3\pi^2} \int_\Delta^\infty \frac{u \, du}{\sqrt{u^2 - \Delta^2}} (u - \frac{q}{2})^3 \theta(\frac{q}{2} - \Delta)
\]

\[
= \frac{N_c}{3(4\pi)^2} \left[ q \sqrt{\frac{q^2}{4} - \Delta^2} (26\Delta^2 + q^2) - 12\Delta^2 (\Delta^2 + q^2) \log \frac{q + 2\sqrt{\frac{q^2}{4} - \Delta^2}}{2\Delta} \right] \theta(\frac{q}{2} - \Delta)
\]

\[
= f(\Delta, q).
\]

We next integrate \( V_{\text{div}, \pm} \) using dimensional regularization. This is done by first integrating over \( p_\perp \) and then over \( u \). This yields

\[
V_{\text{div}} = V_{\text{div}, +} + V_{\text{div}, -} = 2N_c \left( \frac{e^\gamma A^2}{\Delta^2} \right)^\epsilon \left[ 2\Delta^4 \Gamma(-2 + \epsilon) + q^2 \Delta^3 \Gamma(\epsilon) + \frac{q^4}{12} (-1 + \epsilon) \Gamma(1 + \epsilon) \right].
\]

Expanding \( V_{\text{div}} \) to zeroth order in \( \epsilon \), we obtain

\[
V_{\text{div}} = 2N_c \left( \frac{e^\gamma A^2}{\Delta^2} \right)^\epsilon \left[ \left( \frac{1}{\epsilon} + \frac{3}{2} \right) \Delta^4 + \frac{1}{\epsilon} \Delta^2 q^2 - \frac{q^4}{12} + O(\epsilon) \right].
\]

The one-loop effective potential is then given by the sum of Eqs. (63) and (65). It contains poles in \( \epsilon \), which are removed by mass and coupling-constant renormalization. This amounts to making the substitutions \( m^2 \to Z_m m^2 \), \( \lambda \to Z_\lambda \lambda \), and \( g^2 \to Z_{g^2} g^2 \), where

\[
Z_m = 1 + \frac{4N_c g^2}{(4\pi)^2 \epsilon}, \quad Z_\lambda = 1 + \frac{8N_c}{(4\pi)^2 \epsilon} [\lambda g^2 - 6g^4], \quad Z_{g^2} = 1 + \frac{4N_c g^2}{(4\pi)^2 \epsilon}.
\]

In contrast to calculation in the 1+1 dimensional NJL model, we were not able to calculate directly in dimensional regularization the vacuum energy given by \( V_1 = V_+ + V_- \). We therefore use another strategy. In order to isolate the ultraviolet divergences, we expand the integrand in powers of \( q \) and identify appropriate subtraction terms. This yields
After renormalization, the effective potential in the mean-field approximation reads

\[ V = \frac{1}{2} q^2 \Delta^2 + \frac{1}{2} m^2 \Delta^2 \frac{\Lambda^2}{g^2} + \frac{\lambda}{24} \Delta^4 + \frac{2N_c q^2 \Delta^2}{(4\pi)^2} \log \frac{\Lambda^2}{\Delta^2} + \frac{2N_c \Delta^4}{(4\pi)^2} \left[ \log \frac{\Lambda^2}{\Delta^2} + \frac{3}{2} \right] - \frac{N_c q^4}{6(4\pi)^2} + f(\Delta, q). \]  

(67)

In contrast to the example in 1+1 dimensions, we need not subtract a term proportional to the appropriate power of the wavevector (here \( q^4 \)) to obtain an effective potential with the right properties. The reason is simply that the vacuum energy is independent of \( q \) for \( \Delta = 0 \).

We close this section by discussing how dimensional regularization can be used in conjunction with a Landau-Ginzburg (GL) analysis of the quark-meson model. In this case we expand the effective potential in powers of \( \Delta \) and its derivatives. Up to a temperature-dependent constant, we find

\[ V = \frac{1}{2} q^2 \Delta^2 + \frac{1}{2} m^2 \Delta^2 \frac{\Lambda^2}{g^2} + \frac{\lambda}{24} \Delta^4 + \beta_3 \Delta^2 \]

\[ + \beta_2 \Delta^4 + \beta_3 (\nabla \Delta)^2 + \ldots , \]  

(68)

where the coefficients are

\[ \beta_1 = -4N_c \sum_{\{P\}} \frac{1}{P^2}, \]

\[ \beta_2 = 2N_c \sum_{\{P\}} \frac{1}{P^4}, \]

\[ \beta_3 = -N_c \sum_{\{P\}} \left[ \frac{4P_0^2}{P^6} - \frac{3}{P^4} \right]. \]  

(71)

Here, the sum-integral is defined by

\[ \sum_{\{P\}} = \left( \frac{e^{\gamma_E} \Lambda^2}{4\pi} \right)^{\frac{d}{2}} T \sum_{\{P_0\}} \int \frac{d^d p}{(2\pi)^d}, \]  

(72)

where \( P_0 = (2n + 1)\pi T + i\mu \) are the fermionic Matsubara frequencies with \( n = 0, \pm 1, \pm 2\ldots \). Using integrating by parts in \( d = 3 - 2\epsilon \) dimension, it is straightforward to show that \( \beta_2 = \beta_3 \). This result was first obtained in [10] using Pauli-Villars regularization. For the special value of the sigma mass \( m_\sigma = 2\Delta \), it was shown in [11] that this implies that the tricritical point is actually a Lifschitz point. In the NJL model this is always the case when using a regulator where the total derivative vanishes [13]. Due to infinite surface terms, such an expansion is problematic in the case of a momentum cutoff. This problem is avoided in the NJL model in 1+1 dimensions, since the coefficients in the GL functional are finite.

IV. SUMMARY AND DISCUSSION

In this paper, we have for the first time discussed momentum cutoff regularization, symmetric energy cutoff regularization, and dimensional regularization in the context of one-dimensional inhomogeneities in the NJL and QM models. We have shown that all regularization schemes can be used to define a physically meaningful vacuum energy. In the case of symmetric energy cutoff regularization, the result is independent of the wavevector when the magnitude of the condensate vanishes, while in the other cases one must subtract a wavevector-dependent term. We propose to subtract such a term for all regularizations as a part of the renormalization procedure. In the examples considered in this paper, an appropriate term is the Hamiltonian of a free Fermi gas after a unitary transformation. After this subtraction, one must also add a term that depends on the isospin chemical potential in order to obtain the correct expression for the free energy and isospin density in the limit \( b = 0 \).

We have also briefly discussed finite temperature and a Ginzburg-Landau analysis of critical points. Due to the absence of surface terms in the coefficients of the GL functional, dimensional regularization can always be used in the analysis of critical points. The application of momentum cutoff or symmetric energy cutoff at finite temperature is restricted to the cases where the GL coefficients are finite, for example the NJL model in 1+1 dimensions. Results for the phase diagram of the 1+1 dimensional NJL model is presented in [13].

There are other regularization schemes that we have briefly mentioned, namely Schwinger’s proper time regularization and Pauli-Villars regularization. The latter method was successfully applied to the problem of inhomogeneous phases in the NJL model [13] and the QM model [14], where the equality of the two coefficients \( \beta_2 \) and \( \beta_3 \) was shown. In other words, Pauli-Villars regularization has the same virtues as dimensional regularization although the final expressions for renormalized quantities are not so compact.

It is often argued that since the NJL model in three dimensions is “nonrenormalizable”, one cannot use dimensional regularization but is forced to use cutoff (momentum or energy) regularization or Pauli-Villars regularization. We disagree with this.
Nonrenormalizability alone cannot be an argument against applying dimensional regularization since it has been applied successfully to nonrenormalizable models. For example, it has been used in chiral perturbation theory and in the theory of weakly interacting Bose gases and Bose condensation, both involving nonrenormalizable field theories.

ACKNOWLEDGMENTS

The authors would like to thank Tomas Brauner, Michael Buballa and Bernd-Jochen Schaefer for valuable discussions. The authors would like to thank the Niels Bohr International Academy for its hospitality during the latter stages of this work. P.A. would like to thank Professor Sujeev Wickramasekera (Grinnell College) for his mentorship and inspiring scientific work.

[1] P. Fulde and R. A. Ferrell, Phys. Rev. 135, A550, (1964).
[2] A. Larkin and Y. Ovchinnikov, Zh. Eksp. Teor. Fiz. 47, 1136 (1964).
[3] A. W. Overhauser, Phys. Rev. Lett. 4, 415 (1960).
[4] A. B. Migdal, Rev. Mod. Phys. 50, 107 (1978).
[5] K. Maeda, T. Hatsuda, and G. Baym, Phys. Rev. A 87, 021604 (2013).
[6] M. G. Alford, J. A. Bowers, and K. Rajagopal, Phys. Rev. D 63, 074016 (2001).
[7] R. Anglani, G. Nardulli, M. Ruggieri, and M. Mannarelli, Phys. Rev. D 74, 074005 (2006).
[8] L. He, M. Jin, and P.-F. Zhuang, Phys. Rev. D 75, 036003 (2007).
[9] T. Kojo, Y. Hidaka, L. McLerran, and R. D. Pisarski, Nucl. Phys. A 843, 37 (2010); ibid 875, 94, (2012).
[10] T. Kojo, R. D. Pisarski, and A.M. Tsvelik, Phys. Rev. D 82, 074015 (2010).
[11] M. Sadzikowski and W. Broniowski, Phys. Lett. B 688, 83 (2010).
[12] E. Nakano and T. Tatsumi, Phys. Rev. D 71, 114006 (2005).
[13] D. Nickel, Phys. Rev. D 80, 074025 (2009).
[14] D. Nickel, Phys. Rev. Lett. 103, 072301 (2009).
[15] S. Carignano, D. Nickel, and M. Buballa Phys. Rev. D 82, 054009 (2010).
[16] S. Carignano, M. Buballa, and B.-J. Schaefer, Phys. Rev. D 90, 014033 (2014).
[17] H. Abuki, Phys. Lett. B 728, 427 (2014).
[18] J. Braun, S. Finkbeiner, F. Karbstein, and D. Roscher Phys. Rev. D 91,116006 (2015).
[19] T.-G. Lee, E. Nakano, Y. Tsue, T. Tatsumi, B. Friman, Phys. Rev. D 92, 034024 (2015).
[20] A. Heinz, F. Giacosa, and D. H. Rischke, Nucl.Phys. A 933, 34 (2015).
[21] M. Buballa and S. Carignano, Eur. Phys. J. A 52, 57 (2016).
[22] A. Heinz, F. Giacosa, M. Wagner, and D. H. Rischke, Phys. Rev. D 93, 014007 (2016).
[23] R. Anglani, R. Casalbuoni, M. Ciminale, N. Ippolito, R. Gatto, M. Mannarelli, and M. Ruggieri, Rev. Mod. Phys. 86, 509 (2014).
[24] M. Buballa and S. Carignano, Prog. Part. Nucl. Phys. 81, 39 (2015).
[25] O. Schnetz, M. Thies, and K. Urlichs, Annals Phys. 321, 2604 (2006).
[26] M. Thies, J. Phys. A 39, 12707 (2006).
[27] C. Boehmer, M. Thies, and K. Urlichs, Phys. Rev. D 75, 105017 (2007).
[28] G. Basar, G. V. Dunne, and M. Thies, Phys. Rev. D 79, 105012 (2009).
[29] G. Basar and G. V. Dunne, Phys. Rev. Lett. 100, 200404 (2008); Phys. Rev. D 78, 065022 (2008).
[30] D. Ebert, N. V. Gubina, K. G. Klimenko, S. G. Kurbanov, and V. Ch. Zhukovsky, Phys. Rev. D 84, 025004 (2011).
[31] V. Ch. Zhukovsky, K. G. Klimenko, and I. E. Frolov, Moscow Univ. Phys. Bull. 65, 539 (2010).
[32] N. V. Gubina, K. G. Klimenko, S.G. Kurbanov, and V.Ch. Zhukovsky Phys.Rev. D 86, 085011 (2012).
[33] M. Thies, e-Print: arXiv:1603.06218.
[34] S. P. Klevansky, Rev. Mod. Phys. 64, 649 (1992).
[35] G. Basar, G. V. Dunne, and M. Thies, Phys. Rev. D 82, 034029 (2010).
[36] G. P. Lepage, http://arxiv.org/abs/hep-ph/0505233.
[37] N. D. Mermin and H. Wagner, Phys. Rev. Lett. 17, 1133 (1966).
[38] S. Coleman, Commun. Math. Phys. 31, 259 (1973).
[39] J. O. Andersen and T. Brauner, Phys. Rev. D 81, 096004 (2010).
[40] G. Aarts, D. Ahrensmeier, R. Baier, J. Berges, and J. Serreau, Phys. Rev. D 66, 045008 (2002).
[41] J. O. Andersen, Phys. Rev. D 75, 065011 (2007).
[42] E. S. Scavenius, A. Mocsy, I. N. Mishustin, D. H. Rischke, Phys. Rev. C 64, 045202 (2001).
[43] S. Carignano, D. Nickel, and M. Buballa Phys. Rev. D 92, 034024 (2015).
[44] P. Adhikari and J. O. Andersen, e-Print: arXiv:1610.01647 [hep-th].
[46] J. Gasser and H. Leutwyler, Annals Phys. 158, 142 (1984).
[47] E. Braaten and A. Nieto Eur. Phys. Journal B 11, 143 (1999).
[48] J. O. Andersen, Rev. Mod. Phys. 76, 599 (2004).