

# Mesonic dynamics and QCD phase transition

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We study the finite temperature and density two flavor quark-meson model under the functional renormalisation group. The effect of broken  $O(4)$ -symmetry of the wave function renormalisation and expansion point of effective potential on the thermodynamic quantities and baryon number fluctuation are investigate. At the same time, the field dependent Yukawa coupling is also considered. We give results of the pion mass, the quark mass, the trace anomaly, the baryon number fluctuation and the freeze-out curve.

## I. INTRODUCTION

The QCD phase structure and the search of the critical end point (CEP) are the most popular research direction in both experimental and theoretical field. The phase transition between the quark gluon plasma (QGP) and hadron is the main research objects. The research of the QGP-hadron phase transition can help us to help us better understand the nature of elementary particles. The experiment to looking for the QGP is being made at the Large Hadron Collider (LHC) and the Relativistic Heavy-Ion Collider (RHIC).

In terms of theoretical research, there many different methods to investigate the QCD phase structure. The most widely studied method is the lattice QCD. A lot of properties of the QCD matter have been discussed under the lattice simulations. Although the lattice theory has the sign problem at high baryon chemical potential, it still gave us plenty of great outcomes. In order to avoid the problem that occurs in lattice calculation, the study of the continuous non-perturbative field theory is in progress at the same time. For example, the Dyson-Schwinger equation. And the Functional Renormalization Group (FRG) is the other good functional approach of the continuous theory. In these ways we can study the behavior of the strong interaction matter under the finite temperature and density better.

This work is done with the low energy effective model under the FRG approach.

The low-energy effective models, e.g. the quark-meson (QM) model [1], Nambu–Jona-Lasinio (NJL) model, and their Polyakov-loop improved variants: PQM and PNJL, are suitable to be employed to study the QCD phase transitions. They have been investigated quite a lot in literatures, see, e.g., [2] for more details. In this work, we adopt the scale-dependent effective action for the two-

flavor PQM model, as follows

$$\Gamma_k = \int_x \left\{ Z_{q,k} \bar{q} \left[ \gamma_\mu \partial_\mu - \gamma_0 (\mu + i g A_0) \right] q + \frac{1}{2} \left[ Z_{\phi,k}^{\parallel} (\partial_0 \phi)^2 + Z_{\phi,k}^{\perp} (\partial_i \phi)^2 \right] + h_k(\rho) \bar{q} (T^0 \sigma + i \gamma_5 \vec{T} \cdot \vec{\pi}) q + V_k(\rho) - c \sigma \right\}, \quad (1)$$

with  $\mu = (0, 1, \dots, 3)$  and  $i = (1, 2, 3)$ . In Eq. (1) we have used notation  $\int_x = \int_0^{1/T} dx_0 \int d^3x$ , where the imaginary time formalism for the field theory at finite temperature is used, and the temporal length reads  $\beta = 1/T$ . Apparently, when the temperature is nonzero, the  $O(4)$ -symmetry in the Euclidean space is broken into that of  $\mathbb{Z}_2 \otimes O(3)$ , which leads to the split of the magnetic and electric components of correlations functions. They correspond to the components transversal and longitudinal to the heat bath, respectively. In this work, we take this split into account in the two-point correlation function for the mesons, as shown in the second line on the r.h.s. of Eq. (1), where  $Z_{\phi,k}^{\parallel}$  and  $Z_{\phi,k}^{\perp}$  indicate the longitudinal and transversal wave function renormalizations for the temporal and spacial components, respectively.

The reason why we concentrate on the split of the wave function renormalization especially for the mesons is due to the facts as follows. Firstly, in comparison to the quark wave function renormalization  $Z_{q,k}$  and the scale dependent Yukawa coupling  $h_k$  in Eq. (1), it is found that the meson wave function renormalization  $Z_{\phi,k}$  plays the most significant role beyond the local potential approximation (LPA) [2, 3]. In the LPA, the propagators are classical, i.e.,  $Z_{q,k} = Z_{\phi,k} = 1$  and the Yukawa coupling  $h$  is a constant and independent of the scale  $k$ . Secondly, In Ref. [4] calculations based on the full momentum-dependent two-point correlation functions of mesons are compared with those from LPA and LPA', and here in LPA' a momentum-independent  $Z_{\phi,k}$  is included, and it is found that there is a good agreement between the full momentum calculation and the LPA', while not LPA, which indicates that the dispersion relation for the meson, resulting from a scale dependent  $Z_{\phi,k}$ , have already captured most momentum dependence of the two-point correlation function.

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Considering the importance of the wave function renormalization for the mesons and the success of LPA', in this work we would like to investigate the effects of the splitting of  $Z_{\phi,k}$  in the LPA' as shown in Eq. (1), which is a natural choice at finite temperature as discussed above. Furthermore, we will also study the interplay between the splitting of  $Z_{\phi,k}$  and other truncation approaches, e.g., the field dependent Yukawa coupling  $h_k(\rho)$  which encodes higher order quark-meson scattering processes [2], fixed point expansion for the effective potential  $V_k(\rho)$  in Eq. (1) versus the physical point expansion, etc. Their influences on the QCD phase transition and observables, e.g. fluctuations of the baryon number, will be investigated in detail.

## II. FUNCTIONAL RENORMALIZATION GROUP AND FLOW EQUATIONS

To proceed, we describe other notations in the effective action in Eq. (1).  $\phi = (\sigma, \vec{\pi})$  is a meson field with four components. The effective potential  $V_k(\rho)$  with  $\rho = \phi^2/2$  is  $O(4)$  invariant and the c-term  $c\sigma$  breaks the chiral symmetry explicitly. The mesons interact with quarks through the scalar and pseudo-scalar channels with a mesonic field dependent Yukawa coupling  $h_k(\rho)$ , and  $T^0$  and  $T^i$  are the generators in the flavor space with the convention as follows:  $\text{Tr}(T^i T^j) = \frac{1}{2}\delta^{ij}$  and  $T^0 = \frac{1}{\sqrt{2N_f}}\mathbb{1}_{N_f \times N_f}$  with  $N_f = 2$ . Besides the wave function renormalization for the meson, we also introduce one for the quark, i.e.,  $Z_{q,k}$ . Since it plays a minor role in the chiral dynamics in comparison to  $Z_{\phi,k}$ , the splitting of  $Z_{q,k}$  into the transversal and longitudinal components are neglected for simplicity in calculations. Finally,  $\mu$  in the first line on the r.h.s. of Eq. (1) denotes the quark chemical potential, and here the temporal gluon background field  $A_0$ , which encodes the Polyakov dynamics, are also taken into account.

The renormalization group (RG) scale  $k$  in Eq. (1) is an infrared cutoff, below which quantum fluctuations are suppressed in the effective action. The evolution of the effective action with  $k$  is described by the Wetterich equation [5], which reads

$$\partial_t \Gamma_k[\Phi] = \frac{1}{2} \text{Tr}(G_{\phi\phi,k} \partial_t R_k^\phi) - \text{Tr}(G_{q\bar{q},k} \partial_t R_k^q), \quad (2)$$

with the RG time  $t = \ln(k/\Lambda)$  and the initial ultraviolet (UV) cutoff  $\Lambda$ , where  $R_k^\phi$  and  $R_k^q$  are the regulators for the meson and quark fields, respectively, and they are given in Eqs. (A1) and (A2). The scale dependent meson and quark propagators are given by

$$G_{\phi\phi/q\bar{q}}[\Phi] = \left( \frac{1}{\frac{\delta^2 \Gamma_k[\Phi]}{\delta \Phi^2} + R_k^\Phi} \right)_{\phi\phi/q\bar{q}}, \quad (3)$$

with  $\Phi = (q, \bar{q}, \phi)$  denoting all species of fields.

Inserting the effective action in Eq. (1) into the flow equation (2), one arrives at the flow equation for the effective potential, which reads

$$\begin{aligned} \partial_t V_k(\rho) = & \frac{k^4}{4\pi^2} \left[ (N_f^2 - 1) l_0^{(B)}(\bar{m}_{\pi,k}^2, \eta_{\phi,k}^\perp, z_\phi; T) \right. \\ & + l_0^{(B)}(\bar{m}_{\sigma,k}^2, \eta_{\phi,k}^\perp, z_\phi; T) \\ & \left. - 4N_c N_f l_0^{(F)}(\bar{m}_{q,k}^2, \eta_{q,k}; T, \mu) \right], \quad (4) \end{aligned}$$

with the RG invariant dimensionless meson and quark masses as follows

$$\bar{m}_{\pi,k}^2 = \frac{V'_k(\rho)}{k^2 Z_{\phi,k}^\perp}, \quad \bar{m}_{\sigma,k}^2 = \frac{V'_k(\rho) + 2\rho V''_k(\rho)}{k^2 Z_{\phi,k}^\perp}, \quad (5)$$

$$\bar{m}_{q,k}^2 = \frac{h_k^2 \rho}{2k^2 Z_{q,k}^2}. \quad (6)$$

The threshold functions  $l_0^{(B)}$  and  $l_0^{(F)}$  are presented in ??, and the anomalous dimensions for the mesons and quark are defined as

$$\eta_{\phi,k}^\perp = -\frac{\partial_t Z_{\phi,k}^\perp}{Z_{\phi,k}^\perp}, \quad \eta_{\phi,k}^\parallel = -\frac{\partial_t Z_{\phi,k}^\parallel}{Z_{\phi,k}^\parallel}, \quad \eta_{q,k} = -\frac{\partial_t Z_{q,k}}{Z_{q,k}}. \quad (7)$$

Note that  $z_\phi \equiv Z_{\phi,k}^\parallel / Z_{\phi,k}^\perp$  enters into the flow of  $V_k(\rho)$  through the mesonic fluctuations as shown in Eq. (4).

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## III. POLYAKOV-QUARK-MESON MODEL UNDER THE FRG APPROXIMATION

As is shown in the previous work, the Polyakov-quark-meson model of two-flavor can be studied with the flow equation of the scale-dependent effective action  $\Gamma_k[\Phi]$ ,  $\Phi$  is the superfield, it can be written as  $\Phi = (q, \bar{q}, \phi, \dots)$ . The flow equation within the framework of FRG can be written as

$$\partial_t \Gamma_k[\Phi] = \frac{1}{2} \text{Tr}(G_{\phi\phi} \partial_t R_k^\phi) - \text{Tr}(G_{q\bar{q}} \partial_t R_k^q) \quad (8)$$

$$G_{\phi\phi/q\bar{q}}[\Phi] = \left( \frac{1}{\frac{\delta^2 \Gamma_k[\Phi]}{\delta \Phi^2} + R_k^\Phi} \right)_{\phi\phi/q\bar{q}} \quad (9)$$

is the propagator which is full field-dependent. The effective action that depends on the scale of the quark-meson model can be written like this

$$\begin{aligned} \Gamma_k = \int_x \{ & Z_{q,k} \bar{q} [\gamma_\mu \partial_\mu - \gamma_0 (\mu + i g A_0)] q + \frac{1}{2} Z_{\phi,k} (\partial_\mu \phi)^2 \\ & + h_k \bar{q} (T^0 \sigma + i \gamma_5 \vec{T} \cdot \vec{\pi}) q + V_k(\rho) - c\sigma \} + \dots \quad (10) \end{aligned}$$

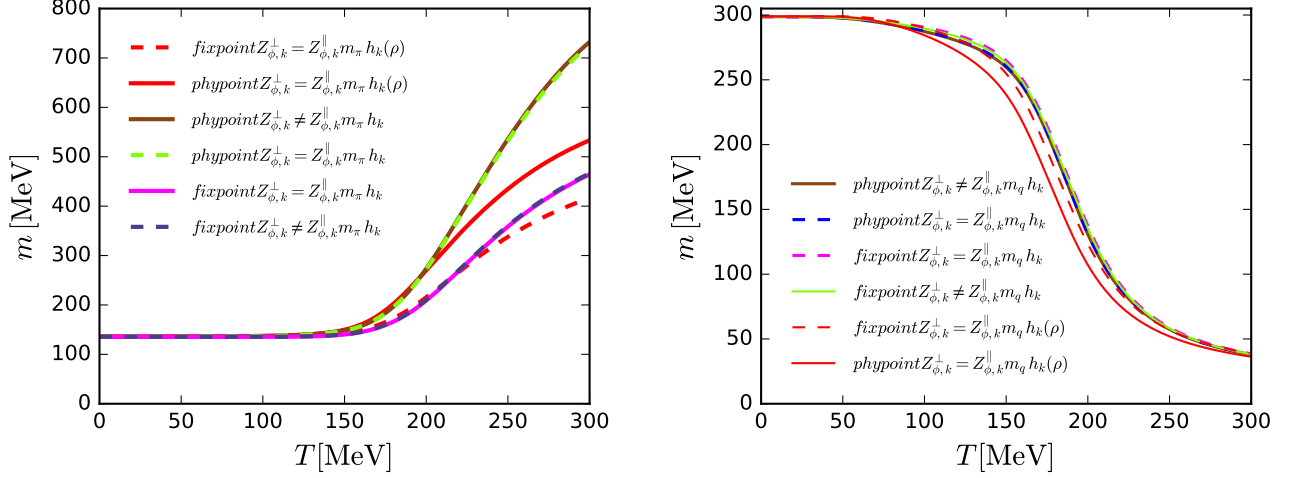


FIG. 1. The left diagram is the mass of the pion as a function of temperature, the right curves are the constituent quark mass. These diagrams are obtained under fix point and physical point expansion of the effective potential.

here we omitted the higher-order terms. The integral sign can be written as  $\int_x = \int_0^{1/T} dx_0 \int d^3x$ . The meson field is  $\phi = (\sigma, \vec{\pi})$ .  $V_k(\rho)$  is meson field-dependent effective potential which is  $O(4)$  invariant, with  $\rho = \phi^2/2$ . The  $k$  is the infrared cutoff scale in FRG;  $\Lambda$  is the UV cutoff renormalization scale, in this work  $\Lambda$  takes 700MeV;  $\mu$  is the chemical potential of quark.  $T$  is the generators of flavor space  $\text{Tr}(T^i T^j) = \frac{1}{2} \delta^{ij}$  and  $T^0 = \frac{1}{\sqrt{2N_f}} \mathbb{1}_{N_f \times N_f}$ . In this work, we only consider the situation of two flavors that  $N_f = 2$ . The linear term  $-c\sigma$  breaks the chiral symmetry. At the same time, the linear breaking parameter  $c$  is related to the mass of  $\vec{\pi}$ ;  $h_k$  is the Yukawa coupling.

In the effective action and the flow equation, we only consider the matter part that is composed of quark loop and meson loop. The glue part is working as a input background. For the purpose to encode the information of de-confinement, we adopted the Polyakov loop. The Polyakov loop plays the role of the background field potential which has the same effect as the gluon. It works in the model by the expectation value of the traced Polyakov loop which can be written as

$$L = \frac{1}{N_c} \langle \text{Tr } \mathcal{P} \rangle, \quad \bar{L} = \frac{1}{N_c} \langle \text{Tr } \mathcal{P}^\dagger \rangle \quad (11)$$

and

$$\mathcal{P}(\vec{x}) = \mathcal{P} \exp \left( ig \int_0^\beta d\tau A_0(\vec{x}, \tau) \right) \quad (12)$$

$L$  and  $\bar{L}$  can be considered as the order-parameter of the phase transition from confinement to the de-confinement.

The bosonic and the fermionic distribution satisfy the

equation given below

$$n_B(\bar{m}_{\phi,k}^2, z_\phi; T) = \frac{1}{\exp\{\frac{1}{T} \frac{k}{z_\phi^{1/2}} \sqrt{1 + \bar{m}_{\phi,k}^2}\} - 1} \quad (13)$$

and

$$n_F(\bar{m}_{q,k}^2, z_q; T) = \frac{1}{\exp\{\frac{1}{T} \frac{k}{z_q} \sqrt{1 + \bar{m}_{q,k}^2} - \mu\} + 1} \quad (14)$$

With the addition of the Polyakov loop, the distribution function of the fermion can be modified by the following form

$$n_F(x, z_q, T, L, \bar{L}) = \frac{1 + 2\bar{L}e^{x/T} + Le^{2x/T}}{1 + 3\bar{L}e^{x/T} + 3Le^{2x/T} + e^{3x/T}} \quad (15)$$

The  $x$  stands for

$$\begin{aligned} x &= \frac{k}{z_q} \sqrt{1 + \bar{m}_{q,k}^2} - \mu \\ \bar{x} &= \frac{k}{z_q} \sqrt{1 + \bar{m}_{q,k}^2} + \mu \end{aligned} \quad (16)$$

The plus and minus sign in front of the chemical potential  $\mu$  stand for the quark and anti-quark.

#### A. The flow equations of the effective potential and the Yukawa coupling

We use the three-dimension regulators throughout our calculation of the threshold functions. Through the derivation of the effective action 8 and 9 we can get the flow equation of the effective potential under the constant mesonic fields:

$$\begin{aligned} \partial_t V_k(\rho) = & \frac{k^4}{4\pi^2} [(N_f^2 - 1) l_0^{(B,4)} (\bar{m}_{\pi,k}^2, \eta_{\phi,k}^\perp; T) \\ & + l_0^{(B,4)} (\bar{m}_{\sigma,k}^2, \eta_{\phi,k}^\perp; T) \\ & - 4N_c N_f l_0^{(F,4)} (\bar{m}_{q,k}^2, \eta_{q,k}; T, \mu)] \end{aligned} \quad (17)$$

The  $l_0^{(B/F,n)}$  stands for the threshold functions of the boson and fermion. The analytical form of the threshold functions are given at ???. Below are the quark mass and meson masses which are renormalized and dimensionless

$$\begin{aligned} \bar{m}_{\pi,k}^2 &= \frac{V'_k(\rho)}{k^2 Z_{\phi,k}^\perp} \\ \bar{m}_{\sigma,k}^2 &= \frac{V'_k(\rho) + 2\rho V''_k(\rho)}{k^2 Z_{\phi,k}^\perp} \\ \bar{m}_{q,k}^2 &= \frac{h_k^2 \rho}{2k^2 Z_{q,k}^2} \end{aligned} \quad (18)$$

In the finite temperature the meson wave function renormalization  $Z_{\phi,k}$  is split into  $Z^\parallel$  and  $Z^\perp$ . In the past calculation we use a approximation that we assume  $Z^\parallel = Z^\perp$ . In this work, we abolished this approximation and observe if the change have any influence on the results. By considering the difference between  $Z_{\phi,k}^\perp$  and  $Z_{\phi,k}^\parallel$  we can calculate the anomalous dimensions respectively. The definition of the anomalous dimensions is given by

$$\eta_{\phi,k}^\perp = -\frac{\partial_t Z_{\phi,k}^\perp}{Z_{\phi,k}^\perp}, \quad \eta_{\phi,k}^\parallel = -\frac{\partial_t Z_{\phi,k}^\parallel}{Z_{\phi,k}^\parallel} \quad (19)$$

and the definition of the quark anomalous dimension is

$$\eta_{q,k} = -\frac{\partial_t Z_{q,k}}{Z_{q,k}} \quad (20)$$

The frequency and spatial momentum are independent when we calculating the anomalous dimensions. So we can easily obtain the longitudinal component of the meson anomalous dimension with the same derivation method of the transversal component. The analytic form of the transversal component is consistent with the anomalous dimension under  $Z_{\phi,k}^\perp = Z_{\phi,k}^\parallel$ . The calculation of anomalous dimensions is ultimately the calculation of the flow of the wave-function renormalizations. We can use the Wetterich equation which is mentioned in Eq 9.

Here we consider two kinds of expansion point of the effective potential. On one hand, the Taylor expansion of the effective potential is about a field value  $\kappa$  which is unrenormalised and fixed. The effective potential can be written as

$$\bar{V}_k(\bar{\rho}) = \sum_{n=0}^{N_v} \frac{\bar{\lambda}_{n,k}}{n!} (\bar{\rho} - \bar{\kappa}_k)^n \quad (21)$$

with  $\bar{\lambda}_{n,k} = \lambda_{n,k}/Z_{\phi,k}^{\perp n}$  and  $\bar{\kappa}_k = Z_{\phi,k}^\perp \kappa$ . The other hand, we make the expansion point  $\kappa$  depends on the cutoff scales  $k$ . Then the  $\bar{\kappa}_k$  can be written as

$$\bar{\kappa}_k = Z_{\phi,k}^\perp \kappa_k \quad (22)$$

Then we give the flow equation of the derivation of the effective potential

$$\begin{aligned} & \partial_{\bar{\rho}}^n (\partial_t|_{\bar{\rho}} \bar{V}_k(\bar{\rho})) \Big|_{\bar{\rho}=\bar{\kappa}_k} \\ &= (\partial_t \bar{\lambda}_{n,k} - n\eta_{\phi,k} \bar{\lambda}_{n,k}) - (\partial_t \bar{\kappa}_k + \eta_{\phi,k} \bar{\kappa}_k) \bar{\lambda}_{n+1,k} \end{aligned} \quad (23)$$

The definition of the pion decay constant is  $f_\pi = \langle \sigma \rangle$  which is equal to the expected value of the  $\sigma$  field.

The Yukawa coupling  $h_k$  is the coupling coefficient of the quark and meson. Here we also consider the  $\rho$ -dependent  $h$  and calculate its impact on the decouple speed of the meson mass. Similar to the expansion of the effective potential, the Yukawa coupling can be expanded at the point  $\kappa_k$ . So the  $h_k$  can be written as

$$\bar{h}_k(\bar{\rho}) = \sum_{n=0}^{N_h} \frac{\bar{h}_{n,k}}{n!} (\bar{\rho} - \bar{\kappa}_k)^n \quad (24)$$

We only study the impact of the expansion of  $h$  under the  $Z_{\phi,k}^\parallel = Z_{\phi,k}^\perp$ , so the definition of the  $\bar{h}_k(\bar{\rho})$  and the expansion coefficients is

$$\bar{h}_k(\bar{\rho}) = \frac{h_k(\rho)}{Z_{q,k} Z_{\phi,k}^{1/2}} \quad (25)$$

$$\bar{h}_{n,k} = \frac{h_{n,k}}{Z_{q,k} Z_{\phi,k}^{(2n+1)/2}} \quad (26)$$

The expression of the Yukawa coupling is

$$\begin{aligned} \partial_t \bar{h}_k(\bar{\rho}) = & \left( \frac{1}{2} \eta_{\phi,k} + \eta_{q,k} \right) \bar{h}_k(\bar{\rho}) \\ & + 8v_3 \bar{h}_k^3(\bar{\rho}) \left[ L_{(1,1)}^{(4)} (\bar{m}_{q,k}^2, \bar{m}_{\sigma,k}^2, \eta_{q,k}, \eta_{\phi,k}; T, \mu) \right. \\ & \left. - (N_f^2 - 1) L_{(1,1)}^{(4)} (\bar{m}_{q,k}^2, \bar{m}_{\pi,k}^2, \eta_{q,k}, \eta_{\phi,k}; T, \mu) \right] \\ & + 4v_3 \bar{h}_k(\bar{\rho}) \bar{h}'_k(\bar{\rho}) \bar{\rho} \left[ \bar{h}_k(\bar{\rho}) + \bar{\rho} \bar{h}'_k(\bar{\rho}) \right] \\ & \times L_{(1,1)}^{(4)} (\bar{m}_{q,k}^2, \bar{m}_{\sigma,k}^2, \eta_{q,k}, \eta_{\phi,k}; T, \mu) \\ & - 2v_3 k^2 \left[ (3\bar{h}'_k(\bar{\rho}) + 2\bar{\rho} \bar{h}''_k(\bar{\rho})) l_1^{(B,4)} (\bar{m}_{\sigma,k}^2, \eta_{\phi,k}; T) \right. \\ & \left. + 3\bar{h}'_k(\bar{\rho}) l_1^{(B,4)} (\bar{m}_{\pi,k}^2, \eta_{\phi,k}; T) \right] \end{aligned} \quad (27)$$

Similar to the effective potential, the flow equation of the derivative Yukawa coupling is given by

$$\begin{aligned} & \partial_{\bar{\rho}}^n (\partial_t|_{\bar{\rho}} \bar{h}_k(\bar{\rho})) \Big|_{\bar{\rho}=\bar{\kappa}_k} \\ &= (\partial_t \bar{h}_{n,k} - n\eta_{\phi,k} \bar{h}_{n,k}) - (\partial_t \bar{\kappa}_k + \eta_{\phi,k} \bar{\kappa}_k) \bar{h}_{n+1,k} \end{aligned} \quad (28)$$

## B. Baryon number fluctuation and kurtosis

The fluctuation of the baryon number is related to some observables. The kurtosis of the baryon number fluctuation at different collision energy can be observed in the experiment like STAR. So the calculation of the baryon number fluctuation is important work. The definition of the each order baryon number fluctuation can be written as

$$\chi_n^B = \frac{\partial^n}{\partial(\mu_B/T)^n} \frac{p}{T^4} \quad (29)$$

The baryon chemical potential satisfying the relationship of  $\mu_B = 3\mu$ . The quadratic and quartic fluctuations can be expressed using the average baryon number fluctuation  $\langle N_B \rangle$

$$\begin{aligned} \chi_2^B &= \frac{1}{VT^3} \langle \delta N_B^2 \rangle \\ \chi_4^B &= \frac{1}{VT^3} (\langle \delta N_B^4 \rangle - 3\langle \delta N_B^2 \rangle^2) \end{aligned} \quad (30)$$

The definition of the kurtosis can be written as

$$\kappa\sigma^2 = \frac{\chi_4^B}{\chi_2^B} \quad (31)$$

At the same time, we calculated the chemical freeze-out line as the function of the collision energy and the kurtosis of the baryon number fluctuation. The value of the kurtosis is determined by the temperature and the baryon chemical potential. We use the correspondence of the collision energy and chemical potential which is give in the Table I. We obtain the freeze-out temperature by the three and two order of the baryon number fluctuation  $\chi_3^B/\chi_2^B$  under the different baryon chemical potential which are corresponding to the different collision energy. Then we can compare the numerical results between the fix point expansion and physical point expansion of the effective potential  $V_k(\rho)$  and the results of splitted meson wave function renormalisation. There are many ways to determine the freeze-out temperature.

$\sqrt{s}$ (GeV)	200	62.4	39	27	19.6	11.5	7.7
$\mu_{B,N_f=2}$ (MeV)	25.3	78.1	121	168.7	222.7	343	459.4

TABLE I. The relationship of collision energy and baryon chemical potential

## IV. NUMERICAL RESULTS AND CONCLUSION

In this work, we investigated the baryon number fluctuations and some thermodynamic quantities with in the functional renormalization group. We made a comparison of the results between the approximation of the mesonic wave function renormalizations and without the

approximation. The calculation are accomplished under the fix point and physical point expansion of the effective potential. Now we give the initial conditions of our calculation. The initial UV scale is  $\Lambda = 700\text{MeV}$ . At  $k = \Lambda$  the initial effective potential can be written as

$$V_\Lambda(\rho) = \frac{\lambda_\Lambda}{2}\rho^2 + \nu_\Lambda\rho \quad (32)$$

And the parameter we chosen in the equation above is shown in the table below. These four parameters are

	$\lambda_\Lambda$	$\nu_\Lambda$	c	h
$\kappa Z_\phi^\parallel = Z_\phi^\perp h_k$	10.15	8.11( $\text{GeV}^2$ )	0.2568( $\text{GeV}^3$ )	7.274
$\kappa Z_\phi^\parallel \neq Z_\phi^\perp h_k$	10.15	8.17( $\text{GeV}^2$ )	0.2568( $\text{GeV}^3$ )	7.274
$\kappa_k Z_\phi^\parallel = Z_\phi^\perp h_k$	5.00	17.86( $\text{GeV}^2$ )	0.3696( $\text{GeV}^3$ )	10.110
$\kappa_k Z_\phi^\parallel \neq Z_\phi^\perp h_k$	11.00	17.91( $\text{GeV}^2$ )	0.3684( $\text{GeV}^3$ )	10.198

TABLE II. The setting of the initial condition parameters.

chosen by fitting the observables:  $m_\pi = 135.9\text{MeV}$ ,  $m_\sigma = 460.0\text{MeV}$ ,  $f_\pi = 92.1\text{MeV}$ ,  $m_q = 298.8\text{MeV}$ .

In the end, we give the numerical calculation results of the physical quantity. In this section we will show the results of the comparison in the two flavor PQM model. Two comparisons have been made, one is  $Z_\phi^\parallel = Z_\phi^\perp$  and  $Z_\phi^\parallel \neq Z_\phi^\perp$ , the other one is the fix point and physical point expansion of the effective potential.

As we can see in the Fig. II, the pion mass and constituent quark mass as the function of temperature. It is clear that the pion mass under the physical point is larger at the high temperature, which means the pion is decoupling quicker. For the meson mass is related to the expansion of the effective potential, the expansion point is running with the cutoff scale  $k$  or not will of course influence the value of the pion mass. However, the split of the meson wave function renormalisation make little effect on the mass.

As is shown in the Fig. III B, because the baryon chemical potential has little influence to the numerical result of the  $\chi_2^B$  and  $\chi_4^B$ , however the  $\chi_1^B$  and  $\chi_3^B$  are more sensitive to the chemical potential. Consequently, we give the odd order of the fluctuation under the chemical potential of 50 MeV and 100 MeV, and the even order under vanishing chemical potential. The results of the quartic fluctuations under two kinds of wave function renormalizations have larger difference at the peak of the curves, the results of  $Z_{\phi,k}^\perp \neq Z_{\phi,k}^\parallel$  are little lower than the results of  $Z_{\phi,k}^\perp = Z_{\phi,k}^\parallel$ .

In the Fig. 3 we show  $\kappa\sigma^2 = \chi_4^B/\chi_2^B$  the kurtosis of the baryon number fluctuation. We can tell from the curves, the two kinds of  $Z_\phi$  almost no effect on the kurtosis results. On the other hand, the influence of the running expansion point of the effective potential is greater. The physical point expansion suppress the kurtosis quicker as the temperature rises than the fix point expansion. It is

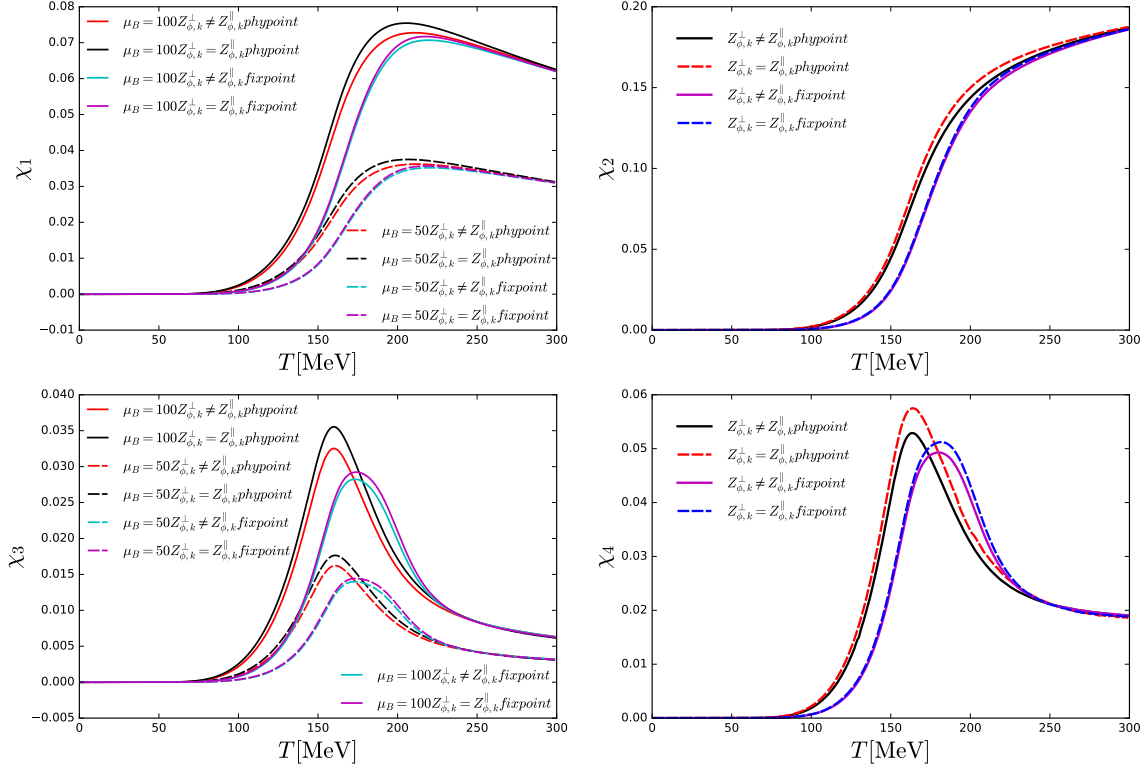


FIG. 2. The first to the fourth order of the baryon number fluctuation, which are obtained under  $Z_\phi^\parallel = Z_\phi^\perp$ ,  $Z_\phi^\parallel \neq Z_\phi^\perp$  and fix point, physical point expansion of the effective potential.

expected from the results of the four order of the fluctuation.

In the Fig. IV we give the pion decay constant and trace anomaly as a function of temperature. The left one is the pion decay constant. It is clear, the separation of the meson wave function renormalisation doesn't make much change. However, the running expansion point of the effective potential brings some differences at high temperature. The trace anomaly doesn't change much either under the different wave function renormalisation, and the physical point expansion suppress the value at high temperature.

## ACKNOWLEDGMENTS

Thanks

## Appendix A: Some useful formulas

In this work we use the  $3d$ - flat or Litim regulators [6, 7], which are very suited for the computations at finite temperature and densities, since the summation for

the Matsubara frequencies is not affected by the  $3d$  regulators, and can be performed analytically. The regulators in Eq. (2) read

$$R_k^\phi(q_0, \mathbf{q}) = Z_{\phi,k}^\perp \mathbf{q}^2 r_B(\mathbf{q}^2/k^2), \quad (\text{A1})$$

$$R_k^q(q_0, \mathbf{q}) = Z_{q,k} i\gamma \cdot \mathbf{q} r_F(\mathbf{q}^2/k^2), \quad (\text{A2})$$

with

$$r_B(x) = \left( \frac{1}{x} - 1 \right) \Theta(1-x), \quad (\text{A3})$$

$$r_F(x) = \left( \frac{1}{\sqrt{x}} - 1 \right) \Theta(1-x), \quad (\text{A4})$$

where  $\Theta(x)$  is the Heaviside step function. Note that since  $Z_{\phi,k}^\perp \neq Z_{\phi,k}^\parallel$ , it is the transversal wave function renormalization for the meson appearing in Eq. (A1).

in which Then we give a definition to the meson and

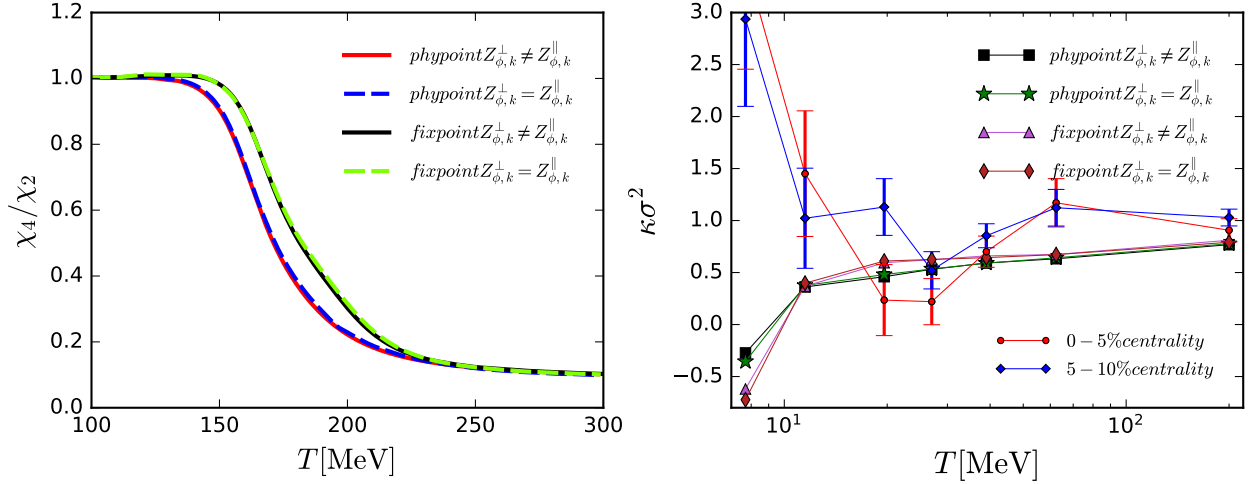


FIG. 3. The left curves are the kurtosis as the function of temperature. The right broken lines are the kurtosis as the function of collision energy.

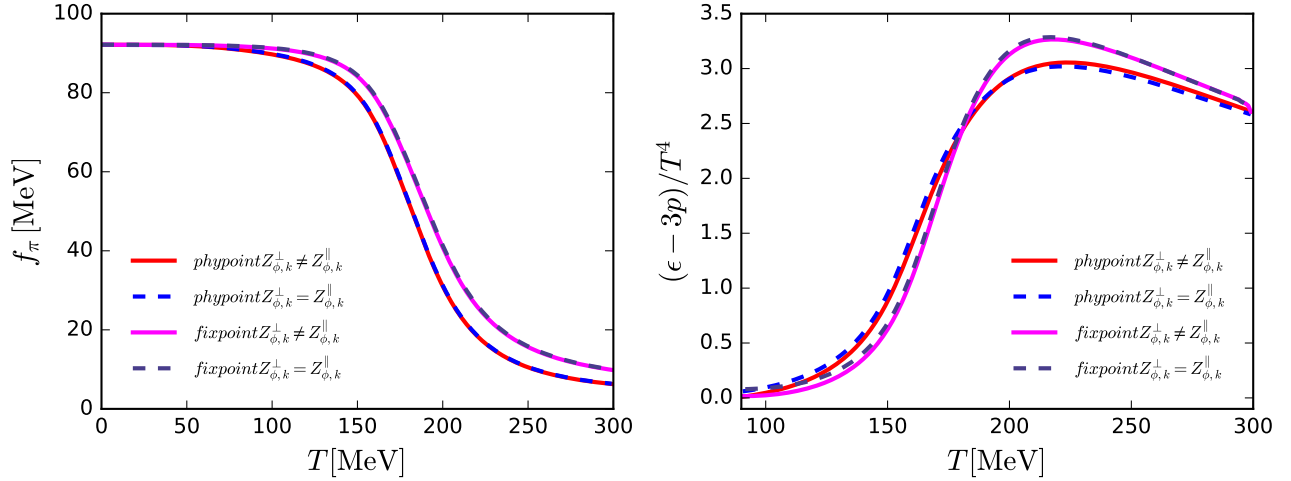


FIG. 4.

quark propagator

$$G_\phi(q, \bar{m}_{\phi,k}^2) = \frac{1}{z_\phi \tilde{q}_0^2 + 1 + \bar{m}_{\phi,k}^2}, \quad (A5)$$

$$G_q(q, \bar{m}_{q,k}^2) = \frac{1}{z_q^2 (\tilde{q}_0 + i\tilde{\mu})^2 + 1 + \bar{m}_{q,k}^2}$$

in the equation above we have  $\tilde{q}_0 = q_0/k$ ,  $\tilde{\mu} = \mu/k$  and for the fermions we have  $q_0 = (2n_q + 1)\pi T$  ( $n_q \in \mathbb{Z}$ ) for the bosons we have  $q_0 = 2n_q\pi T$ . Here  $z = Z^\parallel/Z^\perp$  is the ratio of the two components of the wave function renormalizations. In this work we choose  $z_q = 1$  which means  $Z_q^\parallel = Z_q^\perp$  and  $z_\phi \neq 1$  which means  $Z_\phi^\parallel \neq Z_\phi^\perp$ . To obtain the threshold functions, we define

in the equations above the form of the distribution

functions are

$$n_B(\bar{m}_{\phi,k}^2, z_\phi; T) = \frac{1}{\exp\{\frac{1}{T} \frac{k}{z_\phi^{1/2}} \sqrt{1 + \bar{m}_{\phi,k}^2}\} - 1} \quad (A6)$$

and

$$n_F(\bar{m}_{q,k}^2, z_q; T) = \frac{1}{\exp\{\frac{1}{T} \frac{k}{z_q} \sqrt{1 + \bar{m}_{q,k}^2} - \mu\} + 1} \quad (A7)$$

In the effective potential's flow equation, there are bosonic and fermionic threshold functions. The anomalous dimension of the mesonic field has change into the transverse component of it. The form of the threshold

functions are

$$l_0^{(B,d)}(\bar{m}_{\phi,k}^2, \eta_{\phi,k}^\perp; T) = \frac{2}{d-1} \left( 1 - \frac{\eta_{\phi,k}^\perp}{d+1} \right) \mathcal{B}_{(1)}(\bar{m}_{\phi,k}^2, z_\phi; T) \quad (\text{A8})$$

and

$$l_0^{(F,d)}(\bar{m}_{q,k}^2, \eta_{q,k}; T, \mu) = \frac{2}{d-1} \left( 1 - \frac{\eta_{q,k}}{d} \right) \mathcal{F}_{(1)}(\bar{m}_{q,k}^2, z_q = 1; T, \mu) \quad (\text{A9})$$

The definition of the threshold function  $\mathcal{BB}_{(1,1)}$  is

$$\mathcal{BB}_{(1,1)}(\bar{m}_{\phi_a,k}^2, \bar{m}_{\phi_b,k}^2, z_\phi; T) = -\frac{T}{k} \sum_{n_q} G_\phi(q, \bar{m}_{\phi_a,k}^2) G_\phi(q, \bar{m}_{\phi_b,k}^2) \quad (\text{A10})$$

then we can obtain the  $\mathcal{BB}_{(2,2)}$  in a same way

$$\mathcal{BB}_{(2,2)}(\bar{m}_{\phi_a,k}^2, \bar{m}_{\phi_b,k}^2, z_\phi; T) = \frac{\partial^2}{\partial \bar{m}_{\phi_a,k}^2 \partial \bar{m}_{\phi_b,k}^2} \mathcal{BB}_{(1,1)}(\bar{m}_{\phi_a,k}^2, \bar{m}_{\phi_b,k}^2, z_\phi; T) \quad (\text{A11})$$

The expression of the  $\mathcal{BB}_{(1,1)}$  is

$$\begin{aligned} \mathcal{BB}_{(1,1)}(\bar{m}_{\phi_a,k}^2, \bar{m}_{\phi_b,k}^2, z_\phi; T) = & -\frac{1}{z_\phi^{1/2}} \left\{ \left( \frac{1}{2} + n_B(\bar{m}_{\phi_a,k}^2, z_\phi; T) \right) \frac{1}{(1 + \bar{m}_{\phi_a,k}^2)^{1/2}} \right. \\ & \times \frac{1}{\bar{m}_{\phi_a,k}^2 - \bar{m}_{\phi_b,k}^2} + \left( \frac{1}{2} + n_B(\bar{m}_{\phi_b,k}^2, z_\phi; T) \right) \\ & \left. \times \frac{1}{(1 + \bar{m}_{\phi_b,k}^2)^{1/2}} \frac{1}{\bar{m}_{\phi_b,k}^2 - \bar{m}_{\phi_a,k}^2} \right\} \end{aligned} \quad (\text{A12})$$

then we can get the expression of the threshold functions of any  $n$ . The definition of the  $L_{(1,1)}^{(4)}$  in the flow equation of the Yukawa coupling is

$$\begin{aligned} L_{(1,1)}^{(4)}(\bar{m}_{q,k}^2, \bar{m}_{\phi,k}^2, \eta_{q,k}, \eta_{\phi,k}; T, \mu) = & \frac{2}{3} \left[ \left( 1 - \frac{\eta_{\phi,k}}{5} \right) \mathcal{FB}_{(1,2)}(\bar{m}_{q,k}^2, \bar{m}_{\phi,k}^2; T, \mu) \right. \\ & \left. + \left( 1 - \frac{\eta_{q,k}}{4} \right) \mathcal{FB}_{(2,1)}(\bar{m}_{q,k}^2, \bar{m}_{\phi,k}^2; T, \mu) \right] \end{aligned} \quad (\text{A13})$$

The higher order of the function  $L$  can be obtain by the derivation of the square of mass. The expression of the threshold function  $\mathcal{FB}$  is

$$\begin{aligned} \mathcal{FB}_{(1,1)}(\bar{m}_{q,k}^2, \bar{m}_{\phi,k}^2; T, \mu, p_0) = & \frac{T}{k} \sum_{n_q} G_\phi(p - q, \bar{m}_{\phi,k}^2) G_q(q, \bar{m}_{q,k}^2) \end{aligned} \quad (\text{A14})$$

## Appendix B

In order to obtain the flow equation of the effective potential, the mesonic anomalous dimensions are needed. Because we have divided the wave function renormalizations into the transversal and longitudinal components, so the anomalous dimensions should be divided either. The analytical form of the mesonic anomalous dimensions can be written like

$$\begin{aligned} \eta_{\phi,k}^\parallel = & \frac{1}{6\pi^2} \left\{ \frac{4}{k^2 z_\phi^4} \bar{\kappa}_k (\bar{V}_k''(\bar{\kappa}_k))^2 \left[ -6\mathcal{BB}_{(2,2)}(\bar{m}_{\pi,k}^2, \bar{m}_{\sigma,k}^2; T) \right. \right. \\ & + \frac{4}{z_\phi} (1 + \bar{m}_\sigma^2) \mathcal{BB}_{(2,3)}(\bar{m}_{\pi,k}^2, \bar{m}_{\sigma,k}^2; T) \\ & + \left. \frac{4}{z_\phi} (1 + \bar{m}_\pi^2) \mathcal{BB}_{(3,2)}(\bar{m}_{\pi,k}^2, \bar{m}_{\sigma,k}^2; T) \right] \\ & \left. \times \left( 1 - \frac{1}{5} \eta_{\phi,k}^\perp + \frac{N_c \bar{h}_k^2}{z_\phi} \mathcal{F}_{(3)}(\bar{m}_{q,k}^2; T, \mu) (4 - \eta_{q,k}) \right) \right\} \end{aligned} \quad (\text{B1})$$

and the other component

$$\begin{aligned} \eta_{\phi,k}^\perp = & \frac{1}{6\pi^2} \left\{ \frac{4}{k^2 z_\phi^4} \bar{\kappa}_k (\bar{V}_k''(\bar{\kappa}_k))^2 \mathcal{BB}_{(2,2)}(\bar{m}_{\pi,k}^2, \bar{m}_{\sigma,k}^2; T) \right. \\ & + N_c \bar{h}_k^2 \left[ \mathcal{F}_{(2)}(\bar{m}_{q,k}^2; T, \mu) (2\eta_{q,k} - 3) \right. \\ & \left. \left. - 4(\eta_{q,k} - 2) \mathcal{F}_{(3)}(\bar{m}_{q,k}^2; T, \mu) \right] \right\} \end{aligned} \quad (\text{B2})$$

Now we can obtain the form of the quark anomalous dimension

$$\begin{aligned} \eta_{q,k} = & \frac{1}{24\pi^2 N_f} (4 - \eta_{\phi,k}) \bar{h}_k^2 \\ & \times \left\{ (N_f^2 - 1) \mathcal{FB}_{(1,2)}(\bar{m}_{q,k}^2, \bar{m}_{\pi,k}^2; T, \mu, p_{0,ex}) \right. \\ & \left. + \mathcal{FB}_{(1,2)}(\bar{m}_{q,k}^2, \bar{m}_{\sigma,k}^2; T, \mu, p_{0,ex}) \right\} \end{aligned} \quad (\text{B3})$$

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