Confinement in hot gluonic matter with imaginary and real rotation

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Inspired by recent heavy-ion collision experiments and lattice-QCD simulations, we perturbatively compute the Polyakov loop potential at finite imaginary angular velocity and high temperature. This imaginary rotation does not violate the causality and the thermodynamic limit is well defined. We analytically show that the imaginary angular velocity induces the perturbatively confined phase and serves as a new probe to confinement physics. We discuss a possible phase diagram that exhibits adiabatic continuation from this perturbative confinement to the confined phase at low temperature. Finally we mention subtlety in the analytical continuation from imaginary to real angular velocity by imposing a causality bound.

Introduction: Confinement of quarks and gluons in quantum chromodynamics (QCD) has been a long-standing problem in quantum chromodynamics. There are traditional strategies to idealize the problem. One can take special limits such as the strong-coupling limit [1, 2], the large-$N_c$ limit [3], etc. to scrutinize the confinement mechanism in a nonperturbative and yet analytical way. Such deformations of the theory belong to a category which we can call the QCD-like theory approach. Examples in this category include a supersymmetric extension [4–7]. In particular, an interesting idea of the “adiabatic continuation” on a small circle [7–9] based on the Polyakov mechanism [10] has been recognized. Holographic QCD models are also to be regarded as QCD-like theories [11–13].

Another stream of research toward the confinement mechanism is the introduction of external parameters corresponding to extreme environments such as the temperature $T$, the density or the chemical potential $\mu$, the magnetic field $B$, and so on, which we may call the extreme QCD approach [14, 15]. QCD at extremely high temperature is perturbatively tractable and the loop calculation of the confinement order parameter, i.e., the Polyakov loop, has led to a conclusion that the physical state with the minimum energy is located at the perturbative vacuum in the deconfined phase [16–20] (for a review, see Ref. [21]). Generally speaking, extreme environments provide an energy scale greater than the QCD scale, so that the perturbative calculation in favor of the deconfined phase is justified. The perturbative analysis breaks down with decreasing $T/\mu/B$, and it is usually impossible to go directly into the confinement regime. Nevertheless, one may perceive a precursory tendency of transition from the deconfined to the confined phase (see also Refs. [22, 23] for approaches to enforce confinement to come closer to the transition). One could also employ other external probes like the electric field $E$ [24], the isospin chemical potential $\mu_{\text{iso}}$ [25, 26], the scalar curvature $R$ [27, 28], the rotational angular velocity $\omega$ [29–36], and their mixtures [37, 38].

Since an extraordinary value of $\omega \sim 10^{22}$ s$^{-1}$ was reported in the heavy-ion collision experiment [39], especially, the effect of as large $\omega$ as the QCD scale has been attracting theoretical and experimental interests. Model calculations implied similarity between the angular velocity and the chemical potential [30], which was summarized in a form of the QCD phase diagram on a $\omega-T$ plane [31]. The lattice-QCD simulation suffers from the sign problem at finite $\omega$ in the same way as the finite-$\mu$ case. However, the lattice-QCD simulation with the analytical continuation from the imaginary angular velocity $\Omega$ to $\omega$ via $\Omega = -i\omega$ is feasible because $\Omega$ does not cause the sign problem [29, 34, 35]. Here, we shall emphasize that such a system with imaginary rotation $\Omega$ is quite intriguing on its own. With explicit calculations we show that the pure Yang-Mills theory at sufficiently large $\Omega$ goes through a phase transition from the deconfined to the confined phase even perturbatively.

One might also obtain the perturbatively confining phase by adding finely tuned quark contents, such as a massless adjoint Dirac fermion with imaginary chemical potential, $\mu_1 = \pi$. However, the present work is the very first report of perturbative confinement from purely gluonic loops, to the best of our knowledge. In the literature confining mechanisms that are essentially gluonic are all nonperturbative. They hinge on either semiclassical contribution [7–10], lattice regularization [1, 2], electricity-magnetism duality [4–6], or dressed ghost/gluon propagators [40, 41]. In contrast to these preceding works, it is surprising to find that purely gluonic confinement in (3+1)D is possible without invoking any nonperturbative machinery, just with nontrivial geometry.

Polyakov loop potential with imaginary rotation: We perform the one-loop calculation to find the Polyakov loop potential which is often called the Gross-Pisarski-Yaffe-Weiss (GPY-W) potential [16–18]. Treating rotation effects, we adopt the cylindrical coordinates, i.e., $x^\mu = (\tau, \theta, r, z)$, as well as the metric $g_{\mu\nu} = \text{diag}(1, r^2, 1, 1)$. The geometry of imaginary rotation is specified by the following boundary condition:

\begin{equation}
(\tau, \theta, r, z) \sim (\tau + \beta, \theta - \tilde{\Omega}_t, r, z),
\end{equation}
where $\beta = 1/T$ is inverse temperature and $\tilde{\Omega}_1 := \Omega_1/T$. Clearly, $\tilde{\Omega}_1$ and $\tilde{\Omega}_1 + 2\pi$ describe the same geometry. In the presence of the Polyakov loop or $A_\tau$ background, $\partial_\tau$ is replaced by the covariant derivative $D_\tau$ as

$$D_\tau = \partial_\tau + i \frac{\phi}{\beta} \cdot H.$$  \hfill (2)

The $g$-valued vector $H$ is an orthonormal basis of a Cartan subalgebra of $g$, the Lie algebra of the gauge group. Thus the Polyakov loop is labeled with a real vector $\phi$. We take homogeneous $\phi$ backgrounds because they are the classical vacuum even in the presence of $\Omega_1$.

To perform the one-loop integral, we need to diagonalize the fluctuation operator. For ghosts, it is the scalar Laplacian, $-D^2_v = -D^2 - r^{-1} \partial_r (r \partial_r) - r^{-2} \partial_\theta^2 - \partial_z^2$. We solve the eigenequation, $-D^2_s \Phi = \lambda \Phi$, with the twisted boundary condition (1) to find the spectrum. Since we are merely interested in a potential of $\phi$, we drop the eigenmodes that commute with $H$. Then we find,

$$\Phi_{n,m,k,\alpha}(x) = \frac{E_{\alpha}}{\sqrt{2\pi} \beta} e^{i\left(\frac{2n\pi + \Omega_1 m}{T} \tau + m \theta + k z\right)} J_m(k \perp r).$$  \hfill (3)

Here, $n, m \in \mathbb{Z}$, $k := (k_1, k_2) \in \mathbb{R}^+ \times \mathbb{R}$ and $\alpha$'s are positive roots of $g$. The eigenvalues are given by

$$\lambda_{n,m,k,\alpha} = \left(\frac{2\pi n + \phi \cdot \alpha}{\beta} + \Omega_1 m\right)^2 + |k|^2.$$  \hfill (4)

We can generalize the above calculation to the covariant vector fields, for which the Laplacian is a $4 \times 4$ matrix given by

$$-D^2_v = \begin{pmatrix}
-D^2_s & 0 & 0 & 0 \\
0 & -r^{-1}D^2 r^{-1} + r^{-2} & -2r^{-1} \partial_\theta & 0 \\
0 & 2r^{-3} \partial_\theta & -D^2_s + r^{-2} & 0 \\
0 & 0 & 0 & -D^2_s \\
\end{pmatrix}.$$  \hfill (5)

Its eigenvalues are the same as Eq. (4) but its eigenmodes come with a degeneracy of four polarizations. The unphysical polarizations are simply replicas of the scalar mode (3), i.e., $\Xi_{n,m,k,\alpha}^{(i)}(x) = \Phi_{n,m,k,\alpha}(x) \xi_i^{(i)}$, where $\xi^{(1)} := (1, 0, 0, 0)^T$ and $\xi^{(2)} := (0, 0, 0, 1)^T$. The loop of these unphysical eigenmodes are canceled by the ghost loop. The physical transverse eigenmodes have nontrivial tensorial structure with $m$ shifted by the helicity of the vector fields as

$$\Xi_{n,m,k,\alpha}^{(\pm)}(x) = \frac{E_\alpha E_{\alpha}}{2\sqrt{\pi} \beta} e^{i\left(\frac{2n\pi + \Omega_1 m}{T} \tau + m \theta + k z\right)} J_{m\pm1}(k \perp r),$$  \hfill (6)

where $\xi^{(\pm)} := (0, r, \pm i, 0)^T$.

After performing the Matsubara summation and dropping the ultraviolet divergence independent of $\phi$, we find the following expression for the effective potential:

$$V = \frac{T}{4\pi^2} \sum_{\alpha} \sum_{m \in \mathbb{Z}} \int_0^\infty k_\perp dk_\perp \int_{-\infty}^\infty dk_z \left[ J^2_{m-1}(k_\perp r) + J^2_{m+1}(k_\perp r) \right] \text{Re} \ln \left[ 1 - e^{-(|k|^2 - \Omega_1 m)/T + i\phi \cdot \alpha} \right].$$  \hfill (7)

Interestingly, we can analytically perform the summation and integrals using the power series $\ln(1 - z) = -\sum_{i=1}^\infty \frac{z^i}{i}$, which converges for $|z| \leq 1, z \neq 1$. We then obtain a simple expression,

$$V(\phi; \Omega_1) = -\frac{2T^4}{\pi^2} \sum_{\alpha} \sum_{i=1}^\infty \cos(\phi \cdot \alpha) \cos\left(\frac{i\Omega_1}{2\pi}\right).$$  \hfill (8)

where we introduced dimensionless $\tilde{r} := rT$. Clearly, at $\Omega_1 = 0$, Eq. (8) loses its $r$-dependence and recovers the well-known GPY-W potential [16–18].

The $\phi$-vacua predicted by our inhomogeneous potential (8) exhibit very mild $\tilde{r}$-dependence and are, in particular, homogeneous in the vicinity of $\tilde{\Omega}_1 = 0 \mod \pi$. Actually, as we shall see shortly, the most nontrivial physics reported in this Letter exactly inhabits this most reliable region. For a concrete reference, we shall focus on the rotation center, $\tilde{r} = 0$, where we can complete the $\tilde{l}$ summation to find

$$V(\phi; \tilde{\Omega}_1)|_{\tilde{l}=0} = \frac{\pi^2 T^4}{3} \sum_{\alpha = \pm 1} B_4\left(\frac{\phi \cdot \alpha + s\tilde{\Omega}_1}{2\pi}\right)_{\text{mod} 1}.$$  \hfill (9)

Here $B_4(x) = x^4 - 2x^3 + x^2 - \frac{1}{30}$ is the 4th Bernoulli polynomial. The physical contents of Eq. (9) are quite rich despite its simple appearance.

**Phase transitions:** Let us investigate the evolution of the Polyakov loop potential with increasing $\tilde{\Omega}_1$. We shall first focus on the simplest case of the color SU(2) theory, in which we have only one component, $\phi := \phi \cdot \alpha$.

Figure 1 shows the evolution of the Polyakov loop potential in terms of $\phi/2\pi$ with increasing $\Omega_1$ at $\tilde{r} = 0$. The solid curve in Fig. 1 reproduces the center breaking GPY-W potential with minima located at $\phi = 0$ and $2\pi$. The positive curvature around the minima then corresponds to the Debye screening mass that stabilizes the deconfined phase at high temperature [42]. We clearly see that the curvature is suppressed as $\tilde{\Omega}_1$ gets larger, and eventually the sign of the curvature flips around $\tilde{\Omega}_1$. Then, the potential minima deviate from the deconfined vacua and the confined vacuum at $\phi = \pi$ is energetically favored. We can visualize this phase transition by plotting $\langle L \rangle$, the expectation value of the fundamental Polyakov loop $L$, as a function of $\Omega_1$ as shown in Fig. 2. We see that $\langle L \rangle$ starts to decrease since $\tilde{\Omega}_1 = (1 - 1/\sqrt{3})\pi$. The dropping curve hits $\langle L \rangle = 0$ at $\tilde{\Omega}_1 = \pi/\sqrt{3}$, indicating a second-order confinement phase transition.
We can intuitively understand the confining force at \( \tilde{\Omega_I} = \pi \) from the twisted geometry (1). It assigns all odd-\( m \) physical transverse modes (6) the fermionic antiperiodic boundary condition. But these modes still preserve bosonic statistics such that their loops have no -1 overall sign. Such an antiperiodic gluon loop has the same net effect as a periodic gluino loop, and thus reverses the one-loop potential.

We move on to the SU(3) case. The positive roots are \( \alpha_1 = (1,0) \), \( \alpha_2 = (1/2, \sqrt{3}/2) \), and \( \alpha_3 = (1/2, -\sqrt{3}/2) \). Accordingly, the order parameter has two components, namely, \( \phi = (\phi_1, \phi_2) \). Modulo periodicities and the Weyl group, \( \phi \) runs in a triangular region spanned by the vertices \((0,0)\), \((2\pi, 2\pi/\sqrt{3})\), and \((2\pi, -2\pi/\sqrt{3})\), as drawn in Fig. 3. The points in this triangle bijectively represent conjugacy classes of SU(3). The \( \mathbb{Z}_3 \) center symmetry acts on this equilateral triangle as its rotational geometry symmetry.

We show the SU(3) potential height in the form of the contour plot in Fig. 3. The light (dark) color indicates the region of larger (smaller) potential values. The left in Fig. 3 shows the potential profile at \( \tilde{\Omega_I} = 0 \). The minima are located at \((0,0)\) and its center symmetry images, which indicates the spontaneous broken of center symmetry. With increasing \( \tilde{\Omega_I} \), these minima depart from the conventional vacua as we observed in the SU(2) case in Fig. 2. A crucial difference from the SU(2) case is, as shown in the middle of Fig. 3, the center symmetric point \((4\pi/3,0)\) is pushed down and eventually at \( \tilde{\Omega_I} = \pi/2 \) we see degeneracy between three shifted deconfined vacua and the center symmetric point. The degeneracy indicates a first-order phase transition, and the center symmetric (confining) state is energetically favored for \( \tilde{\Omega_I} = \pi \) as shown in the right of Fig. 3. We can also visualize this first-order nature again by plotting \( L \) as shown in Fig. 2. Clearly, we see a sudden jump of \( L \) at \( \Omega_I = \pi/2 \). This difference between the phase transition order of SU(3) and that of SU(2) is consistent with the universality argument.

Our formulae hold for any semisimple Lie algebra. We can show that, for any simply connected compact group with a nontrivial center, Eq. (9) always favors center symmetric vacua at \( \tilde{\Omega_I} = \pi \). For example, SU\((N_c > 3)\) or Spin\( (5) \) also exhibits a first-order confinement phase transition at \( \tilde{\Omega_I} = \pi/2 \). A more interesting case is \( G_2 \) which has no center symmetry. Consistently, we observed no phase transition; the location of its potential minimum just continuously moves as a function of \( \tilde{\Omega_I} \).

The potential curvature at the vacua characterizes the correlation length of the matter. It corresponds to the Debye mass square \( m_D^2 \) in the deconfined phase and the string tension \( \sigma \) in the confined phase, respectively. For example, we can find in the SU(3) case that \( m_D^2(\tilde{\Omega_I} = 0) \simeq g^2T^2 \) and \( \sigma(\tilde{\Omega_I} = \pi) \simeq \frac{1}{4}g^2T^2 \).

Here, we discuss an intriguing question on the relation between the ordinary confined phase at low \( T \) and \( \tilde{\Omega_I} = 0 \) and the perturbatively confined phase we found. It is instructive to consider the phase diagram on the \( \tilde{\Omega_I} - T \) plane of as shown in Fig. 4. Our calculations cannot constrain the phase diagram except for the high-\( T \) region. Nevertheless, we may anticipate that the tendency of less deconfinement at larger \( \tilde{\Omega_I} \) persists around \( T_c \), as shown by an increasing \( T_c(\tilde{\Omega_I}) \) function in Fig. 4. The region at low \( T \) and \( \tilde{\Omega_I} \sim \pi \) is totally inaccessible, but it might be an educated guess that the imaginary rotation also reverse nonperturbative effects. Thus we may conjecture a nonperturbative deconfined phase in
the window of $\pi/2 \lesssim \tilde{\Omega} \lesssim 3\pi/2$ at low $T$. So far, these conjectured features are reasonably motivated, but there are two possibilities about the high-$T$ extension of the nonperturbative deconfined phase. If it stays at low $T$ surrounded by the confined phase as in the left of Fig. 4, the perturbative confined phase is smoothly connected to the ordinary confined phase. Then our perturbative confinement may be a handle to catch some features of nonperturbative confinement. However, the right of Fig. 4 is also possible, in which the two confinement phases do not communicate with each other. We cannot exclude this scenario logically, though it looks artificial.

**Analytical continuation to real rotation:** We finally apply our results to the real rotation. It is customary in the literature to study the real rotation effect by the analytical continuation from $\Omega_I$ to $\omega$ [34, 35]. For example, once $T_c(\Omega_I)$ is known, then $T_c(\omega)$ is inferred from the replacement of $\Omega_I^2 = -\omega^2$. However, we explicate that such a procedure might be problematic using our perturbative expression.

For any unreal complex $\tilde{\Omega}_I$, Eq. (7) always yields a potential with singularity at some $\phi$. In fact, our derivation of Eq. (8) is valid for real $\Omega_I$ only and a nonzero Im $\Omega_I$ would drive the Maclaurin series of $\ln(1 - z)$ out of its convergence radius. If we naively perform the analytical continuation to Eq. (8), we would also encounter problems. For $\tilde{r} > 0$, Eq. (8) is analytical everywhere except on the imaginary $\Omega_I$ axis. There, infinitely many poles are accumulated around $\tilde{\Omega}_I = 0$. As for $\tilde{r} = 0$, the poles are gone, but the infinite summation just blows up for unreal $\tilde{\Omega}_I$.

The physical origin of these singularities is clear. At finite angular velocity $\omega$, the longwave modes with $k_\perp \lesssim \omega$ violate the causality so we should introduce an infrared cutoff. That is, $r\omega$ must not exceed the unity. Let us set the system size as $r \leq R$ with $R\omega \leq 1$. This discretizes the momentum $k_\perp$ such that $k_\perp R$ is a zero of the Bessel functions. Here, we denote the $\kappa$-th nontrivial zero of $J_\nu(\xi_{\nu,\kappa})$ as $\xi_{\nu,\kappa}$. Now, the phase space integration in Eq. (7) is replaced as follows:

$$
\int_0^{\infty} k_\perp dk_\perp J^2_m(k_\perp R) f(k_\perp) \rightarrow \sum_{\kappa=1}^{\infty} \frac{2}{R^{2}J^2_{m+1}(\xi_{m,\kappa})} J^2_m\left(\frac{\xi_{m,\kappa} R}{R}\right) f\left(\frac{\xi_{m,\kappa}}{R}\right).
$$

We have performed the numerical integration and summation of Eq. (7) for real $\omega$ with Eq. (10) substituted. In practice we cut off the sum over $m$, $\kappa$ and the $k_z$ integration by sufficiently large numbers.

Figure 5 shows the evolution of the SU(2) Polyakov loop potential (made dimensionless with $T^4$) for $\omega R = 0, 1/2, 1$ in the color SU(2) case at $\tilde{r} = 0$. Our choice of parameters is: $R = 10$ GeV$^{-1}$ and $T = 0.15$ GeV.

$J_\nu(\xi)$ as $\xi_{\nu,\kappa}$. Now, the phase space integration in Eq. (7) is replaced as follows:

$$
\int_0^{\infty} k_\perp dk_\perp J^2_m(k_\perp R) f(k_\perp) \rightarrow \sum_{\kappa=1}^{\infty} \frac{2}{R^{2}J^2_{m+1}(\xi_{m,\kappa})} J^2_m\left(\frac{\xi_{m,\kappa} R}{R}\right) f\left(\frac{\xi_{m,\kappa}}{R}\right).
$$

We have performed the numerical integration and summation of Eq. (7) for real $\omega$ with Eq. (10) substituted. In practice we cut off the sum over $m$, $\kappa$ and the $k_z$ integration by sufficiently large numbers.

Figure 5 shows the evolution of the SU(2) Polyakov loop potential as a function of $\phi/2\pi$ with increasing $\omega$ at $\tilde{r} = 0$. We chose the system size as $R = 10$ GeV$^{-1} (\approx 2$ fm) and $T = 0.15$ GeV. The potential minima are located at $\phi = 0 \bmod 2\pi$ for any $\omega$, so that the system always stays in the deconfined phase. Yet, we can quantify the effect of $\omega$ onto the vacuum stability by evaluating the potential curvature around the minimum. The
potential curvature represents the Debye screening mass squared. As we see in Fig. 5, the curvature increases with increasing $\omega$, and this means that deconfinement is more favored by rotation. However, these Polyakov loop potentials exhibit a sensitive dependence on the system size $R$ in an unpleasant way. Details about this subtlety will be reported elsewhere. The justification of analytical continuation to real rotation needs further investigations.

Outlook: An intriguing and immediate extension of our work would be the lattice simulation to explore the whole $\Omega_I-T$ phase structure as shown in either Fig. 4. Actually, we can have the lattice simulation at our fingertips for $\Omega_I = \pi/2$ and $\pi$ by moving to the Cartesian coordinates, $(\tau, x, y, z)$, where the boundary condition (1) reduces to $(\tau, x, y, z) \sim (\tau + \beta, y, -x, z)$ at $\Omega_I = \pi/2$ and $(\tau, x, y, z) \sim (\tau + \beta, -x, -y, z)$ at $\Omega_I = \pi$, respectively. Therefore, for $\Omega_I = \pi/2$ and $\pi$, we do not have to deal with nontrivial geometry, but just take the square lattice and the Cartesian spacetime with only modification of a twisted thermal boundary condition. Once we manage to know the physics at $\tilde{\Omega} = \pi/2$ and $\pi$ from low to high temperature, we can justify or falsify our speculated phase diagrams in Fig. 4.

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[1] K. G. Wilson, Phys. Rev. D 10, 2445 (1974).
[2] A. M. Polyakov, Phys. Lett. B 72, 477 (1978).
[3] G. ’t Hooft, Nucl. Phys. B 272, 461 (1974).
[4] N. Seiberg and E. Witten, Nucl. Phys. B 426, 19 (1994), [Erratum: Nucl.Phys.B 430, 485–486 (1994)], arXiv:hep-th/9407087.
[5] N. Seiberg and E. Witten, Nucl. Phys. B 431, 484 (1994), arXiv:hep-th/9408099.
[6] N. Seiberg, Nucl. Phys. B 435, 129 (1995), arXiv:hep-th/9411149.
[7] N. M. Davies, T. J. Hollowood, V. V. Khoze, and M. P. Mattis, Nucl. Phys. B 559, 123 (1999), arXiv:hep-th/9905015.
[8] E. Poppitz, T. Schäfer, and M. Unsal, JHEP 10, 115, arXiv:1205.0290 [hep-th].
[9] K. Aitken, A. Cherman, E. Poppitz, and L. G. Yaffe, Phys. Rev. D 96, 096022 (2017), arXiv:1707.08971 [hep-th].
[10] A. M. Polyakov, Nucl. Phys. B 120, 429 (1977).
[11] J. Babington, J. Erdmenger, N. J. Evans, Z. Guralnik, and I. Kirsch, Phys. Rev. D 69, 066007 (2004), arXiv:hep-th/0306018.
[12] M. Kruczenski, D. Mateos, R. C. Myers, and D. J. Winters, JHEP 05, 041, arXiv:hep-th/0311270.
[13] T. Sakai and S. Sugimoto, Prog. Theor. Phys. 113, 843 (2005), arXiv:hep-th/0412141.
[14] K. Rajagopal and F. Wilczek, The Condensed matter physics of QCD, in *At the frontier of particle physics. Handbook of QCD*. Vol. 1-3, edited by M. Shifman and B. Ioffe (2000) pp. 2061–2151, arXiv:hep-ph/0011333.
[15] K. Fukushima, J. Phys. G 39, 013101 (2012), arXiv:1108.2939 [hep-ph].
[16] N. Weiss, Phys. Rev. D 24, 475 (1981).
[17] N. Weiss, Phys. Rev. D 25, 2667 (1982).
[18] D. J. Gross, R. D. Pisarski, and L. G. Yaffe, Rev. Mod. Phys. 53, 43 (1981).
[19] C. P. Korthals Altes, Nucl. Phys. B 420, 637 (1994), arXiv:hep-th/9310195.
[20] A. Gocksch and R. D. Pisarski, Nucl. Phys. B 402, 657 (1993), arXiv:hep-ph/9302233.
[21] K. Fukushima and V. Skokov, Prog. Part. Nucl. Phys. 96, 154 (2017), arXiv:1705.00718 [hep-ph].
[22] A. Vuorinen and L. G. Yaffe, Phys. Rev. D 74, 025011 (2006), arXiv:hep-ph/0604100.
[23] K. Fukushima and N. Su, Phys. Rev. D 88, 076008 (2013), arXiv:1210.8250 [hep-lat].
[24] D. T. Son and M. A. Stephanov, Phys. Atom. Nucl. 64, 834 (2001), arXiv:hep-ph/0011365.
[25] D. T. Son and M. A. Stephanov, Phys. Rev. Lett. 86, 592 (2001), arXiv:hep-ph/0005225.
[26] T. Inagaki, T. Muta, and S. D. Odintsov, Prog. Theor. Phys. Suppl. 127, 93 (1997), arXiv:hep-th/9711084.
[27] A. Flachi and K. Fukushima, Phys. Rev. Lett. 113, 091102 (2014), arXiv:1406.6548 [hep-th].
[28] A. Yamamoto and Y. Hirono, Phys. Rev. Lett. 111, 081601 (2013), arXiv:1303.6292 [hep-lat].
[29] H.-L. Chen, K. Fukushima, X.-G. Huang, and K. Mamedov, Phys. Rev. D 93, 104052 (2016), arXiv:1512.08974 [hep-ph].
[30] Y. Jiang and J. Liao, Phys. Rev. Lett. 117, 192302 (2016), arXiv:1606.03808 [hep-ph].
[31] M. Chernodub, Phys. Rev. D 103, 054027 (2021), arXiv:2012.04924 [hep-ph].
[32] X. Chen, L. Zhang, D. Li, D. Hou, and M. Huang, JHEP 07, 132, arXiv:2010.14478 [hep-ph].
[33] V. V. Braguta, A. Y. Kotov, D. D. Kuznedelev, and A. A. Roenko, Pisma Zh. Eksp. Teor. Fiz. 112, 9 (2020).
[34] V. V. Braguta, A. Y. Kotov, D. D. Kuznedelev, and A. A. Roenko, Phys. Rev. D 103, 094515 (2021), arXiv:2102.05084 [hep-lat].
[35] Y. Fujimoto, K. Fukushima, and Y. Hidaka, Phys. Lett. B 816, 136184 (2021), arXiv:2101.09173 [hep-ph].
[36] A. Flachi and K. Fukushima, Phys. Rev. D 98, 096011 (2018), arXiv:1702.04753 [hep-th].
[37] H. Zhang, D. Hou, and J. Liao, Chin. Phys. C 44, 111001 (2020), arXiv:1812.11787 [hep-ph].
[38] L. Adamczyk et al. (STAR), Nature 548, 62 (2017), arXiv:1701.06657 [nucl-ex].
[39] D. Zwanziger, Nucl. Phys. B 323, 513 (1989).
[40] J. Braun, H. Gies, and J. M. Pawlowski, Phys. Lett. B 684, 262 (2010), arXiv:0708.2413 [hep-th].
[41] A. K. Rebhan, Nucl. Phys. B 430, 319 (1994), arXiv:hep-ph/9408262.