Stabilizing perturbative Yang-Mills thermodynamics with Gribov quantization

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We evaluate the thermodynamic quantities of Yang-Mills theory using the Gribov quantization, which deals with nonperturbative resummation. The magnetic scale is automatically incorporated into the framework and we find it efficient to stabilize the perturbative expansion of the free energy. In the temperature range \(T = T_c \sim 2T_c\), the major uncertainty in our results comes from the nonperturbative running coupling that is adopted from the lattice simulation, while the convergence above \(2T_c\) is impressively robust. We also present the corresponding interaction measure (i.e., trace anomaly) up to close to \(T_c\).

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I. INTRODUCTION

Stimulated by the exciting developments from the Relativistic Heavy Ion Collider (RHIC) and Large Hadron Collider (LHC) experiments, the thermodynamics of the quark-gluon plasma—especially the pressure \(p\) or the free energy density \(f = -p\)—is of crucial interest. Due to asymptotic freedom \([1]\), perturbation theory was expected to work for QCD thermodynamics at high enough temperatures. The weak-coupling expansion of the QCD free energy has been accomplished up to \(g^6\log g\), which is the highest order in the perturbative approach, with \(g\) being the running coupling (see Ref. \([2]\) for reviews). Unexpectedly, the resulting weak-coupling series shows poor convergence in the intermediate-temperature regime, i.e., \(T = 2T_c \sim 4T_c\), with \(T_c \sim 160\) MeV being the pseudocritical temperature for the QCD phase transition (or \(T_c \sim 270\) MeV for pure Yang-Mills theory, abbreviated as YM hereafter), which is probed in the RHIC and LHC experiments. Resummation must be carried out to incorporate the contributions from the electric scale \(gT\) \([3–6]\).

The nonconvergence of QCD thermodynamics is attributed to the IR sector of QCD, which is governed by soft-scale gluons \([2]\). Therefore the pure Yang-Mills theory provides the simplest test bench to access the IR problem. While the resummed perturbation theory in the electric sector resulted in improved convergence of the YM free energy down to \(T \sim 3T_c\), significant deviations from the lattice data were observed by lowering \(T\) further toward \(T_c\), which could be taken as a signal of the onset of nonperturbative effects.

It was discovered in the early 1980s that due to the absence of an IR cutoff by screening in the magnetic sector, the perturbative expansion of finite-\(T\) YM theory breaks down at a fundamental level at the magnetic scale \(g^2 T\). This IR catastrophe is the so-called Linde problem \([7]\). As a result, perturbation theory works up to a certain order only depending on the quantity in consideration. For the free energy IR divergences from the magnetic scale start entering at four loops, corresponding to \(g^6\) order, which makes the highest accessible order \(g^6 \log g\), as mentioned before. In the RHIC and LHC temperature regime the running coupling is \(g \sim \mathcal{O}(1)\), and therefore the missing contributions from the nonperturbative magnetic scale should be as significant as the ones from the perturbative electric scale. This fact provides us with an explanation for the aforementioned deviations in the free energy of resummed perturbation theory from the lattice data. There have been various attempts to incorporate the magnetic contributions, but the nonperturbative nature is inherent in the confining properties of dimensionally reduced YM theory at high temperature \([8]\). We must thus augment the resummed perturbation theory with a confinement mechanism even when we deal with a deconfined state of matter.

The key ingredients in constructing the YM thermodynamics are the correlation functions of the gluons and the ghosts. Gauge fixing is conveniently done through the Faddeev-Popov (FP) procedure \([9]\), and the bare gluon and ghost propagators show an IR pole proportional to \(1/p^2\), with \(p = (p_0, p)\) being the Minkowskian four-momentum. The corresponding on-shell dispersion relation is \(p_0 = |p|\), which is modified in the full propagators including interaction effects.

In order to have a better understanding in the intermediate-temperature regime, the YM correlation functions must be improved nonperturbatively, including the magnetic sector. The last decade has witnessed tremendous developments of high-precision lattice measurements of the YM correlation functions at both zero and finite temperatures \([10]\). There has also been considerable progress in the study of YM correlation functions by functional methods in, e.g., Ref. \([13]\), and the results are in good agreement with the lattice simulation (see Ref. \([14]\) and references therein). The obtained gluon and ghost propagators show an IR suppression and enhancement, respectively, as compared to the FP case. These results heuristically encompass desirable features of confinement: the IR-suppressed gluon propagator indicates gluon confinement at large distance, and the IR-enhanced ghost is responsible for confinement. It has also been demonstrated that the balance in the IR sector is re-
sponsible for the deconfinement phase transition [11, 12].

Gribov pointed out more than three decades ago that there still remains a residual gauge ambiguity in the IR sector, i.e., the so-called Gribov copy problem [15]. If the functional integration is dominated by the contributions near the horizon of the fundamental modular region, the ghost propagator is naturally enhanced in the IR momenta, and the resulting gluons are suppressed in turn. This is in line with the results from the lattice and functional methods. The corresponding on-shell dispersion relation reads

\[ p_0 = \sqrt{p^2 + m_G^2}/p^2, \]

with \( m_G \) being the Gribov mass parameter which is solved by the condition of the horizon dominance. It was shown by Zwanziger in a phenomenological way that a free gas of Gribov quasiparticles qualitatively captures the nonperturbative features of the lattice equation of state [16]. Subsequently, it was found that the Gribov mass parameter is correctly proportional to the magnetic scale, i.e., \( m_G \sim g^2 T \), in the limit \( T \to \infty \) [17]. This finding is promising enough and a confining mechanism can indeed resolve the Linde problem by providing a nonperturbatively generated IR cutoff. Many of the works on the Gribov quantization are available in the limiting cases as [15] for a review). Systematic attempts to extend the theory of confinement to the finite-\( T \) problems should deserve further investigations along the same line as a first try [17]. In this work we will pursue this possibility to make the perturbative evaluation of YM free energy stabilized by the nonperturbative mass scale emerging from the Gribov gauge-fixing procedure in which features of confinement are implemented, which we simply call the Gribov quantization.

II. FORMALISM

In the Gribov quantization [15], the YM partition function in Euclidean space reads

\[ Z = \int_{\Omega} DA(x) V(\Omega) \delta(\partial \cdot A) \det[-\partial \cdot D(A)] e^{-S_{YM}}, \]

in which the functional integration should be carried out within the Gribov region defined as

\[ \Omega \equiv \{ A : \partial \cdot A = 0, -\partial \cdot D(A) \geq 0 \}, \]

whereby the Landau gauge is chosen. The restriction of the integration to the Gribov region is realized by inserting a function \( V(\Omega) \) into the partition function (1), where

\[ V(\Omega) = \theta[1 - \sigma(0)] = \int_{-i\infty}^{+i\infty + \epsilon} \frac{d\beta}{2\pi i \beta} e^{\beta [1 - \sigma(0)]} \]

represents the no-pole condition. Here, \( 1 - \sigma(P) \) is the inverse of the ghost dressing function \( Z_G(P) \) [see Eq. (11)]. The integration variable \( \beta \) is identified as the Gribov mass parameter \( m_G \) after a redefinition (see Ref. [18] for technical details of the Gribov quantization).

A. Gap equation

The Gribov mass parameter \( m_G \) is a Lagrange multiplier that is determined by the variational principle, leading to the following gap equation:

\[ \int P \frac{1}{P^4 + m_G^2} = \frac{d}{(d-1)N_c g^2}, \]

where the Euclidean four-momentum reads \( P = (p, p_4) \), with \( p_4 = 2n\pi T \), and \( N_c \) is the number of colors. Our calculation is carried out in dimensional regularization with the MS renormalization scheme, in which the sum-integral is defined as

\[ \int P \equiv \left( \frac{e^{-\mu^2}}{4\pi} \right)^\epsilon T \sum_{p_4=2n\pi T} \int \frac{d^{d-1}p}{(2\pi)^{d-1}}, \]

with \( d = 4 - 2\epsilon \) being the spacetime dimensions.

After carrying out the sum-integral and subtracting the UV divergence (see Ref. [19] for details) the gap equation becomes

\[ 1 = \frac{3N_c g^2}{64\pi^2} \left[ \frac{5}{6} - \ln \left( \frac{m_G^2}{\mu^2} \right) \right] + \frac{4}{im_G} \int_0^\infty dp p^2 \left( \frac{n_B(\omega_-)}{\omega_-} - \frac{n_B(\omega_+)}{\omega_+} \right), \]

where \( \omega \equiv \sqrt{p^2 + im_G^2} \), and \( n_B(x) \equiv (e^{x/T} - 1)^{-1} \) is the Bose-Einstein distribution function. Analytical solutions are available in the limiting cases as

\[ m_G = \begin{cases} \mu \exp \left( \frac{5}{12} \frac{32\pi^2}{3N_c g^2} \right) & (T \to 0), \\ \frac{d-1}{d} N_c \frac{g^2 T}{4\pi} & (T \to \infty), \end{cases} \]

in which it is evident that the magnetic scale emerges at high temperature. In what follows we will solve Eq. (6) numerically to derive \( m_G \) as a function of \( T \) for a given renormalization scale \( \mu \) (which will be fixed later).

B. Gluon propagator

Gribov’s gluon propagator in the Landau gauge reads

\[ D_A(P) = \delta^{ab} \frac{P^2}{P^4 + m_G^2} \left( \delta^{\mu\nu} - \frac{P^\mu P^\nu}{P^2} \right), \]

which is regular as \( P \to 0 \), indicating the IR suppression of confined gluons [15]. To proceed to any calculation involving \( D_A(P) \), we need to specify the running coupling as a function of \( T \). Since we are interested in the IR regime, the perturbative running is not quite appropriate and we have to use the full nonperturbative
running. Instead of solving the nonperturbative resummation, we adopt the lattice results at finite $T$ [20]. In Ref. [20] the running coupling is extracted from the large-distance (IR) and short-distance (UV) behaviors of the heavy-quark free energy, leading to $\alpha_{\text{IR}}$ and $\alpha_{\text{UV}}$, respectively.

Interestingly, the lattice-measured coupling can be nicely fitted by a one-parameter form as

$$\alpha_s(T/T_c) \equiv g^2(T/T_c) / 4\pi = \frac{6 \pi}{11N_c \ln[c(T/T_c)]},$$

with $c = 1.43$ for the IR case and $c = 2.97$ for the UV case, as shown together with the lattice data in Fig. 1. In the intermediate-temperature regime, the running coupling has a substantial scheme dependence and we therefore use the IR and UV values to estimate theoretical uncertainties.

C. Ghost propagator

The no-pole condition (i.e., to not go across the Gribov horizon) requires that the IR limit of the ghost propagator should be enhanced, indicating the proximity to the Gribov horizon [15]. The ghost propagator in the Landau gauge thus reads

$$D_c(P) = \delta^{ab} \frac{1}{1 - \sigma(P)} \frac{1}{P^2},$$

where the ghost dressing function is $Z_G \equiv [1 - \sigma(P)]^{-1}$ and $\sigma(P)$ with gluon ladder diagrams taken into account turns out to be

$$\sigma(P) \equiv N_c g^2 \int \frac{1}{Q^4 + m_G^4} \left( \frac{Q^2}{(Q - P)^2} \right)^2 \left( \delta^{\mu\nu} - \frac{Q^\mu Q^\nu}{Q^2} \right).$$

Using the gap equation (4) one can confirm $1 - \sigma(0) = 0$, indicating the IR enhancement of the ghosts. We note that recent lattice data favor the so-called decoupling solution for which the ghost dressing function stays finite at $P \to 0$, which can be incorporated into a refined Gribov-Zwanziger approach [21]. In this study, however, we keep using the original form (shown above) for simplicity. Different behaviors in the deep IR region hardly affect bulk thermodynamics. In order to verify this, we have inserted a mass parameter to modify the deep-IR behavior of the propagators in favor of the decoupling solution, and have found that the resulting pressure decreases at most less than 10% at $T = T_c$, and only a few percent at $T > T_c$.

At $T = 0$ this integration can be performed with dimensional regularization, yielding

$$1 - \sigma(P) = \frac{N_c g_0^2}{128\pi^2} \left[ -5 + \left( 3 - \frac{m_G^4}{P^4} \right) \ln \left( 1 + \frac{P^4}{m_G^4} \right) + \frac{\pi P^2}{m_G^2} + \frac{2 \left( 3 - \frac{P^4}{m_G^2} \right)}{m_G^2 \left( 3 - \frac{P^4}{m_G^2} \right)} \arctan \frac{P^2}{m_G} \right],$$

where $m_G$ is the ghost mass at $T = 0$, and we find that a choice of $\mu = 1.69$ GeV [in view of Eq. (7) with $g = 3.13$ adjusted by hand] can reproduce the $T = 0$ lattice data of $Z_G$ [22, 23]. Amazingly, with the same $\mu$ fixed at $T = 0$, our finite-$T$ numerical results from Eq. (11) [with $g(T)$ from Eq. (9)] are rather insensitive to $T$, which is perfectly consistent with the lattice simulation. We note that in the direct evaluation of Eq. (11) we imposed a three-momentum cutoff $\Lambda = 1.25\mu$ so that the $T = 0$ numerical results are matched with Eq. (12). Figure 2 shows the numerical results at $T = 200$ MeV and $T = 400$ MeV when $g(T)$ runs with $c = 2.97$. If we use $c = 1.43$, for the $T = 200$ MeV case, $cT$ is too close to $T_c$. As long as $T$ is greater than $T_c$, however, the behavior of $Z_G$ is robust and not contaminated by the uncertainty in $g(T)$.

Now that we have confirmed that the resummed ghost dressing function in our approach is almost $T$ in-

![Figure 1](https://via.placeholder.com/150)

**FIG. 1.** Lattice data of IR (upper) and UV (lower) running couplings at finite $T$ from Ref. [20] and the fit results with a choice of $c$ in Eq. (9).

![Figure 2](https://via.placeholder.com/150)

**FIG. 2.** Ghost dressing functions at $T = 200$ MeV (solid curve) and $T = 400$ MeV (dashed curve) as compared to the lattice data from Ref. [22].
dependent, as observed in the lattice simulation, we can safely use Eq. (12) for the finite-$T$ calculation too.

III. QUASIPARTICLE APPROXIMATION

With the gluon and ghost propagators prepared, we can calculate the free energy. In the two-particle irreducible (2PI) formalism the effective action can be expressed in terms of the full propagator $G$ and full self-energy $\Pi$ as

$$\Gamma = \frac{1}{2} \text{tr} \ln G^{-1} - \frac{1}{2} \text{tr} \Pi G + \Gamma_2[G],$$

where $\Gamma_2[G]$ represents the sum of all 2PI diagrams. In a general quasiparticle approximation we keep only the first term, which should be a reasonable estimate; the variational principle to derive $G$, i.e., $\delta \Gamma / \delta G = 0$, can be formulated to minimize the sensitivity to higher-order corrections, which is often called optimized perturbation theory [24]. This means that the first term in Eq. (13) must be dominant as long as $G$ is self-consistently determined, or equivalently, if we use the full $G$ that should coincide with the self-consistently determined $G$, in principle.

A. Gluon contribution

The gluon contribution to the free energy reads

$$\frac{1}{2} \text{tr} \ln D^{-1}_A = \frac{N_c^2 - 1}{2} \sum_P \left[ 3 \ln \frac{P^4 + m_G^4}{P^2} + \ln P^2 \right],$$

which can be evaluated straightforwardly with dimensional regularization, giving

$$\frac{1}{2} \text{tr} \ln D^{-1}_A = (N_c^2 - 1) \left\{ \frac{\pi^2 T^4}{45} + \frac{3 m_G^4}{32 \pi^2} \left( \frac{3}{2} - \ln \frac{m_G^2}{\mu^2} \right) \right\} + 3 T \sum_p \int_0^\infty dp p^2 \ln(1 - e^{-\omega_p/T}).$$

We can then calculate the gluon contribution numerically using the value of $\mu$ fixed by the ghost propagator and $m_G$ as a numerical solution of the gap equation (6).

B. Ghost contribution

The ghost contribution to the free energy reads

$$\text{tr} \ln D^{-1}_c = - (N_c^2 - 1) \sum_P \ln \left\{ P^2 \left[ 1 - \sigma(P) \right] \right\}$$

$$= (N_c^2 - 1) \left\{ \frac{\pi^2 T^4}{45} - \sum_P \ln \left[ 1 - \sigma(P) \right] \right\},$$

with Eq. (12) substituted for $1 - \sigma(P)$, which is justified numerically. It is technically daunting to do the sum-integral of Eq. (16) analytically, and we will resort to the numerical calculation, in which the UV divergence is subtracted at a certain $T$ sufficiently below $T_c$.

IV. RESULTS AND DISCUSSIONS

With the renormalization scale $\mu$ fixed as $\mu = 1.69$ GeV, we can numerically solve the gap equation (6) to obtain $m_G$ as a function of $T$, which we show in Fig. 3.

Solving the gluon contribution (15) and the ghost contribution (16) numerically with $m_G$ as determined in Fig. 3, we obtain the pressure or the free energy of Gribov quasiparticles scaled by the massless Stefan-Boltzmann limit, $-(N_c^2 - 1) \pi^2 T^4/45$ for $N_c = 3$, as shown in Fig. 4. In technical practice, the ghost contribution is naively UV divergent and needs regularization and subtraction. We regularized the integration by inserting a smooth function, $(1 - \tanh((P^2 - P_0^2)/\Delta P^2))/2$, where we chose $P_0 = 100$ GeV and $\Delta P^2/P_0^2 = 0.4$. Because we numerically process the Matsubara sum in a straightforward way, it is necessary to use a sufficiently smooth regularization in order to avoid unphysical cutoff artifacts. Then we extracted the matter part of the ghost contribution by subtracting the pressure at a temperature far below $T_c$ (which is 50 MeV in this work). Of course the final results are insensitive to this choice of the subtraction point, which we have explicitly checked.

To draw Fig. 4, the SU(3) YM lattice data (blue dots) are taken from Ref. [25]. The red band shows the uncertainty arising from the lattice running couplings, with the lower (solid) and upper (dashed) bounds corresponding to the IR and the UV couplings, respectively, as we have seen in Fig. 1.

It is worth noting that our calculations are done...
in a way consistent with the lattice-measured $\alpha_s(T)$, while the perturbative running coupling with an assumed renormalization scale used in conventional resummed perturbation theory is not compatible with the lattice $\alpha_s(T)$ even at very high $T$ (see discussions in Ref. [20] for more details). Due to the ambiguity in defining $\alpha_s(T)$ on the lattice, we see a rather wide band at $T \lesssim 2T_c$. However, the uncertainty in the resulting free energy gets suppressed significantly at $T \gtrsim 2.5T_c$, which is highly nontrivial, especially considering the fact that there is still a big variation between the IR and UV lattice couplings at such temperatures.

Comparing to the most recent estimate from three-loop hard-thermal-loop perturbation theory (HTLpt) [6], as shown by the gray band in Fig. 4, the uncertainty in our free energy is about 35% that of the HTLpt at 2.5$T_c$, and about 15% at 5$T_c$. We would like to point out, furthermore, that our free energy is consistent with the lattice data in the whole displayed temperature range. In particular the lattice data lie in our band even below 3.5$T_c$, where HTLpt suffers from poor convergence, and therefore a sharp rising of the free energy after the deconfinement phase transition—which indicates a nonperturbative release of new degrees of freedom—is realized somehow even without the inclusion of the Polyakov loop. The free energies from the resummation schemes of Refs. [3, 4] have relatively smaller uncertainties than HTLpt, but they are not successful enough to recover the sharp rising behavior below 2$T_c$, where we consider that the resummation in the magnetic sector is crucial.

All the other thermodynamic quantities can be derived from the free energy and in Fig. 5 we show the $T^4$-scaled interaction measure $I$ (which is sometimes called the trace anomaly), defined as $I = \epsilon - 3p = T^5 \frac{d\sigma}{dp}(p/T^4)$, where $\epsilon$ is the energy density given by $\epsilon = Tdp/dT - p$. The interaction measure $I$ vanishes when the theory preserves scale invariance, and therefore it is an important measure for the breaking of scale invariance by interaction effects and the quantum transmutation of the mass scale. Besides, because $I$ involves a derivative with respect to $T$, the interaction measure is obviously more informative than the pressure $p$. In other words, the consistency seen in the pressure—as in Fig. 4—does not guarantee quantitative agreement in view of $I$.

We present our result in Fig. 5, which shows similar behavior as the pressure in Fig. 4: the uncertainty from the lattice $\alpha_s(T)$ gets highly suppressed for increasing values of $T$ and the lattice data lie in our band for almost the entire regime shown in the plot. It is clear that $m_G$ has a sizable effect on the interaction measure at low $T$ ($\lesssim 2T_c$) compared to the three-loop HTLpt result, which indicates the onset of the nonperturbative magnetic scale. Due to the absence of the effects of the Polyakov loop, the peak in the interaction measure—which indicates the onset (for decreasing $T$) of phase-transition physics—is not manifest in our result.

In closing, we would like to stress that the resummation of the magnetic sector is by itself incorporated through the Gribov quantization, though the quasiparticle approximation may look like too simple an approach. In fact, however, the construction of $\sigma(P)$ explicitly involves the higher-order diagrams, and solving the gap equation is nothing but the variational calculation to carry out the nonperturbative resummation effectively. In this sense, it is neither an accident nor magic that our results are consistent with the lattice data, but rather a natural consequence from the proper reformation of the magnetic contributions.
V. CONCLUSIONS AND OUTLOOK

In this paper we have explored a novel and systematic evaluation of the Yang-Mills free energy using the Gribov quantization. The results, thanks to the improvement at the magnetic scale through the self-consistent solution for $m_{\tilde{\alpha}}$, show robust and stable behavior consistent with the lattice data, which is the first self-consistent calculation that significantly surpasses the first attempt in Ref. [17]. We have evaluated the finite-$T$ ghost propagator from the Gribov quantization and shown its insensitivity to $T$, which to our knowledge is the very first semianalytic result that is in surprising agreement with the latest lattice observation as seen in Refs. [22, 23]. All these results evidently manifest the profound importance of the nonperturbative gauge fixing even in the perturbative evaluation of thermodynamics at high temperature.

There are intriguing future extensions. First, the explicit calculation of the magnetic screening mass and the spatial string tension would provide us with a deeper insight into the Linde problem. Second, the Polyakov-loop coupling would enable us to investigate the first-order deconfinement transition [12]. Third, the running coupling from the functional approaches are qualitatively in line with the lattice result and would help in reducing the uncertainty band [26, 27]. Fourth, it would be desirable to systematize the perturbation theory using the local Gribov-Zwanziger action [28]. Last but not least, the application of the Gribov quantization would have a profound impact on more general QCD problems, e.g., the magnetic properties of QCD matter [29], transverse dynamics of real-time evolution [30], and so on.

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[1] D. J. Gross and F. Wilczek, Phys. Rev. Lett. 30, 1343 (1973); H. D. Politzer, Phys. Rev. Lett. 30, 1346 (1973).
[2] J. P. Blaizot, E. Iancu, and A. Rebbhan, hep-ph/0303185; U. Kraemmer and A. Rebbhan, Rep. Prog. Phys. 67, 351 (2004); J. O. Andersen and M. Strickland, Ann. Phys. (N.Y.) 317, 281 (2005); N. Su, Commun. Theor. Phys. 57, 409 (2012).
[3] J. P. Blaizot, E. Iancu, and A. Rebbhan, Phys. Rev. Lett. 83, 2906 (1999); Phys. Rev. D 68, 025011 (2003).
[4] M. Laine and Y. Schröder, Phys. Rev. D 73, 085009 (2006); A. Hietanen, K. Kajantie, M. Laine, K. Rummukainen, and Y. Schröder, Phys. Rev. D 79, 045018 (2009).
[5] J. O. Andersen, E. Braaten, and M. Strickland, Phys. Rev. Lett. 88, 2139 (1999); J. O. Andersen, L. E. Leganger, M. Strickland, and N. Su, Phys. Lett. B 696, 468 (2011).
[6] J. O. Andersen, M. Strickland, and N. Su, Phys. Rev. Lett. 104, 122003 (2010); J. High Energy Phys. 08 (2010) 113.
[7] A. D. Linde, Phys. Lett. B 96, 289 (1980); D. J. Gross, R. D. Pisarski, and L. G. Yaffe, Rev. Mod. Phys. 53, 43 (1981).
[8] A. K. Rebbhan, Phys. Rev. D 48, 3967 (1993); Nucl. Phys. B430, 319 (1994); W. Buchmuller and O. Philipsen, Nucl. Phys. B443, 47 (1995); G. Alexanian and V. P. Nair, Phys. Lett. B 352, 435 (1995); P. B. Arnold and L. G. Yaffe, Phys. Rev. D 52, 7208 (1995); R. Jackiw and S. -Y. Pi, Phys. Lett. B 308, 131 (1993); Phys. Lett. B 403, 297 (1997); J. M. Cornwall, Phys. Rev. D 57, 3694 (1998); F. Eberlein, Phys. Lett. B 439, 130 (1998); D. Bieletekz, K. Lessmeier, O. Philipsen, and Y. Schröder, J. High Energy Phys. 05 (2012) 058.
[9] L. D. Faddeev and V. N. Popov, Phys. Lett. B 25, 29 (1967).
[10] A. Maas, Phys. Rep. 524, 203 (2013).
[11] J. Braun, H. Gies, and J. M. Pawlowski, Phys. Lett. B 684, 262 (2010).
[12] K. Fukushima and K. Kashiwa, Phys. Lett. B 723, 360 (2013).
[13] L. von Smekal, A. Hauck, and R. Alkofer, Phys. Rev. Lett. 79, 3591 (1997).
[14] A. C. Aguilar, D. Binosi, and J. Papavassiliou, Phys. Rev. D 78, 025010 (2008); C. S. Fischer, A. Maas, and J. M. Pawlowski, Ann. Phys. 324, 2408 (2009); D. R. Campagnani and H. Reinhardt, Phys. Rev. D 82, 105021 (2010); L. Fister and J. M. Pawlowski, arXiv:1112.5440; M. Q. Huber and L. von Smekal, J. High Energy Phys. 04 (2013) 149; A. C. Aguilar, D. Ibáñez, and J. Papavassiliou, Phys. Rev. D 87, 114020 (2013).
[15] V. N. Gribov, Nucl. Phys. B139, 1 (1978).
[16] D. Zwanziger, Phys. Rev. Lett. 94, 182301 (2005).
[17] D. Zwanziger, Phys. Rev. D 76, 125014 (2007).
[18] N. Vandersickel and D. Zwanziger, Phys. Rep. 520, 175 (2012).
[19] J. A. Gracey, Phys. Lett. B 632, 282 (2006); 686, 319(E) (2010).
[20] O. Kaczmarek, F. Karsch, F. Zantow, and P. Petreczky, Phys. Rev. D 70, 074505 (2004); 72, 059903(E) (2005).
[21] D. Dudal, J. A. Gracey, S. P. Sorella, N. Vandersickel, and H. Verschelde, Phys. Rev. D 78, 065047 (2008).
[22] A. Sternebeck, E.-M. Igenfritz, M. Müller-Preussker, A. Schiller, and I. L. Bogolubsky, Proc. Sci. LAT (2006) 076 [arXiv:hep-lat/0610035].
[23] A. Cucchieri, A. Maas, and T. Mendes, Phys. Rev. D 75, 076003 (2007); R. Aouane, V. G. Bornyakov, E. M. Igenfritz, V. K. Mitrusushkin, M. Müller-Preussker, and A. Sternebeck, Phys. Rev. D 85, 094501 (2012).
[24] S. Chiku and T. Hatsuda, Phys. Rev. D 58, 076001.
(1998).
[25] S. Borsányi, G. Endrödi, Z. Fodor, S. D. Katz, and K. K. Szabó, J. High Energy Phys. 07 (2012) 056.
[26] J. Braun and H. Gies, Phys. Lett. B 645, 53 (2007).
[27] Jan M. Pawlowski (private communication).
[28] D. Zwanziger, Nucl. Phys. B323, 513 (1989).
[29] T. Kojo, Y. Hidaka, L. McLerran, and R. D. Pisarski, Nucl. Phys. A843, 37 (2010); T. Kojo, Y. Hidaka, K. Fukushima, L. D. McLerran, and R. D. Pisarski, Nucl. Phys. A875, 94 (2012); T. Kojo and N. Su, Phys. Lett. B 720, 192 (2013); B. Feng, E. J. Ferrer, and V. de la Incera, arXiv:1304.0256.
[30] A. Dumitru, Y. Nara, and E. Petreska, Phys. Rev. D 88, 054016 (2013).