RG flow of the Polyakov-loop potential - first status report -

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Abstract.
We study SU(2) Yang-Mills theory at finite temperature in the framework of the functional renormalization group. We concentrate on the effective potential for the Polyakov loop which serves as an order parameter for confinement. In this first status report, we focus on the behaviour of the effective Polyakov-loop potential at high temperatures. In addition to the standard perturbative result, our findings provide information about the “RG improved” backreactions of Polyakov-loop fluctuations on the potential. We demonstrate that these fluctuations establish the convexity of the effective potential.

Keywords: renormalization group, quark deconfinement, gauge theories
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1. INTRODUCTION

An understanding of strongly interacting matter at finite temperature is a prominent problem of contemporary physics that deserves to be analyzed with great effort in view of the current and future experiments at heavy ion colliders. Since the forces between quarks as elementary constituents are governed by a non-Abelian gauge theory, already the understanding of gauge boson dynamics is an important challenge. In this latter case of pure gluodynamics, the expected transition to a deconfined phase can be studied with the aid of the Polyakov loop [1], being the order parameter for this transition:

$$\mathcal{P}(\vec{x}) = \frac{1}{N} \text{Tr}_F \exp \left( i \bar{g} \int_0^\beta A_0^a(\vec{x}, t) t^a dt \right). \quad (1)$$

Here $t^a$ are the generators of $SU(N)$ and $\beta$ denotes the inverse temperature. The subscript $F$ alludes to the fundamental representation. The negative logarithm of the Polyakov-loop expectation value can be interpreted as the free energy of a single static fundamental color source [2]. In this sense, an infinite free energy associated with confinement is indicated as $\langle \mathcal{P} \rangle \to 0$, whereas $\langle \mathcal{P} \rangle \neq 0$ signals deconfinement.

Moreover, $\langle \mathcal{P} \rangle$ measures whether center symmetry, a discrete symmetry of Yang-Mills theory, is realized by the thermodynamic ensemble [2, 3]. Gauge transformations which differ at Euclidean times $x_0 = 0$ and $x_0 = \beta$ by a center element of the gauge group change $\mathcal{P}$ by a phase $e^{i2\pi k/N}$, $k$ integer, but leave the action and the functional integration measure invariant. This implies that a center-symmetric ground state automatically ensures $\langle \mathcal{P} \rangle = 0$, whereas deconfinement $\langle \mathcal{P} \rangle \neq 0$ is related to the breaking of this symmetry.
In fact, lattice simulations have not only collected strong evidence for a second order phase transition in $SU(2)$ Yang-Mills theory \cite{4, 5}, but reveal moreover that the critical exponents agree with those of a 3D $Z(2)$ Ising model \cite{6}. The latter corresponds exactly to the conjectured universality class obtained from the Polyakov-loop criterion \cite{7}.

In recent years, an effective theory consisting of gauge-invariant powers of the Polyakov loop has been developed based on mean-field arguments, see \cite{8} for an overview. Such considerations show a good agreement with lattice data \cite{6}. Moreover, inverse Monte-Carlo techniques have recently facilitated a precise lattice determination of the Polyakov-loop effective action for $SU(2)$ \cite{10}.

A perturbative calculation of the effective potential $V$ for the order parameter was first performed by Weiss \cite{11}. For this, it is convenient to work with the “Polyakov gauge”, which rotates the zeroth component $A_0$ of the gauge field into the Cartan subalgebra of $SU(N)$; furthermore, the condition $\partial_0 A_0 = 0$ is imposed. Focussing solely on the Polyakov loop, it suffices to consider a single scalar degree of freedom $\phi(\vec{x})$, defined by

$$A(\vec{x})_\mu = n^a \delta_{\mu 0} \phi(\vec{x}).$$  \hspace{1cm} (2)

Here $n^a$ denotes a constant unit vector in color space, e.g. $n^a = \delta^{a3}$. Using this gauge, the order parameter is fully determined by $\phi$,

$$\mathcal{P}(\vec{x}) = \cos \left( \frac{\beta \bar{g} \phi(\vec{x})}{2} \right),$$  \hspace{1cm} (3)

and, consequently, the effective potential can purely be expressed in terms of $\phi$, being a compact variable $\beta \bar{g} \phi \in [0, 2\pi]$. To leading order in a derivative expansion, the effective potential yields \cite{11}

$$\beta^4 V(\beta \bar{g} \phi) = -\sum_{n=1}^{\infty} \frac{4}{n^4 \pi^2} \cos(n \beta \bar{g} \phi),$$  \hspace{1cm} (4)

displayed here as a Fourier-cosine series, and depicted on the left panel of Fig. 1. The potential has minima at $\beta \bar{g} \phi = 2\pi n$ and is $Z(2)$-symmetric, i.e., invariant under $\beta \bar{g} \phi \rightarrow 2\pi - \beta \bar{g} \phi$. Furthermore the order parameter is finite for $\beta \bar{g} \phi = 2\pi n$ and therefore the $Z(2)$-symmetry is spontaneously broken and the system is in the deconfined phase. This perturbative result agrees with the expectation that perturbation theory holds for high temperature where the coupling is small. It fails to describe the confinement phase.

In the confined phase, the potential should have its minimum at $\beta \bar{g} \phi = \pi$, implying a vanishing order parameter. A sketch of a possible form of the potential with a finite IR regulator (to circumvent the convexity obstruction, see below) is shown on the right panel of Fig. 1 for illustration.

Various generalizations to Weiss’s result have been worked out within perturbation theory, for instance, the inclusion of a magnetic background field \cite{12} or higher-order derivative expansions \cite{13} to name a few. However, in order to investigate the transition to the confinement phase, reliable nonperturbative tools are
required. We will base our study on the functional (or “Exact”) renormalization
group [14] formulated in terms of a flow equation for the effective action [15].

The paper is organized as follows: In Sect. 2 we briefly review RG flow equations
for Yang-Mills theories and present the flow equation for the effective potential of
the order parameter in a propertime approximation. The flow of the potential is
analyzed in Sect. 3. Conclusions and future directions are discussed in Sect. 4.

2. FLOW EQUATION FOR THE POLYAKOV-LOOP
POTENTIAL

2.1. Exact Renormalization Group

In the flow equation approach, we consider the effective average action \( \Gamma_k \) which
includes all quantum fluctuations with momenta \( |p| > k \), with the scale \( k \) serving
as an infrared (IR) regularization. The boundary condition for the flow equation is
fixed at an ultraviolet (UV) scale \( \Lambda \) in terms of the bare action \( \Gamma_\Lambda \) to be quantized.
Quantum fluctuations are successively integrated out by lowering the scale \( k \). In the
limit \( k \to 0 \), the full quantum effective action \( \Gamma_k \to 0 \), i.e., the generating functional
of the 1PI Green’s functions, is obtained. The flow of \( \Gamma_k \), i.e., the RG trajectory
from the UV scale \( \Lambda \) to the deep IR, is obtained from a functional differential
equation [15].

Flow equations for gauge theories require a careful control of gauge invariance,
because the IR regulator scale \( k \) introduces sources of gauge-symmetry breaking in
addition to standard gauge-fixing terms. Nevertheless, standard gauge symmetry
can be obtained in the physical limit of vanishing regulator scale \( k \to 0 \) by control-
ling the symmetry constraints with the aid of regulator-modified Ward-Takahashi
identities [16]. In this paper, we employ the flow equation with the background-field
method [17, 18] for a simplified control of gauge invariance within our approxima-
tion. We work along the lines of [17, 19] where this technique has been used for
zero-temperature gluodynamics. The flow equation for the effective average action
reads \[17\]
\[
k \partial_k \Gamma_k[A,\bar{A}] \equiv \partial_t \Gamma_k[A,\bar{A}] = \frac{1}{2} \text{Tr} \frac{\partial_t R_k(\Gamma_k^{(2)}[\bar{A},\bar{A}])}{\Gamma_k^{(2)}[A,\bar{A}] + R_k(\Gamma_k^{(2)}[\bar{A},\bar{A}])}. \tag{5}
\]

Here the trace runs over all internal indices including momenta. The classical gauge field is given by \(A\), whereas the background-field is denoted by \(\bar{A}\) (the ghost fields are not displayed for brevity). \(\Gamma_k^{(2)}\) denotes the second functional derivative with respect to the fluctuating fields. The regulator function \(R_k\) implements the IR regularization at the scale \(k\), see below. Inserting the background-field dependent \(\Gamma_k^{(2)}\) into the regulator leads to an adjustment of the regularization to the spectral flow of the fluctuations as discussed in \[19\], implying a potential improvement when it comes to approximations.

The boundary condition for the effective action at the UV scale \(\Lambda\) consists of:
\[
\Gamma_\Lambda[A,\bar{A}] = \Gamma^{cl}[A,\bar{A}] + \Gamma^{gf}_\Lambda[A,\bar{A}] + \Gamma^{gh}_\Lambda[A,\bar{A}], \tag{6}
\]

with the bare Yang-Mills action \(\Gamma^{cl}\). The gauge-fixing and ghost terms are given by
\[
\Gamma^{gf}_\Lambda[A,\bar{A}] = \frac{1}{2\xi_\Lambda} \int d^d x (D_\mu[\bar{A}]a_\mu)^2, \quad \Gamma^{gh}_\Lambda[A,\bar{A}] = -\int d^d x cD_\mu[\bar{A}]D_\mu[A]c, \tag{7}
\]
respectively, with the quantum fluctuations \(a_\mu\) defined by \(a = A - \bar{A}\). The gauge parameter is denoted by \(\xi_k\).

Let us briefly summarize a few properties of the regulator function \(R_k\) that can conveniently be written as
\[
R_k(x) = xr(y), \quad y := \frac{x}{Z_k k^2}, \tag{8}
\]
with \(r(y)\) being a dimensionless regulator shape function of dimensionless argument. Here \(Z_k\) denotes a wave-function renormalization. Note that both \(R_k\) and \(Z_k\) are matrix-valued in field space.

The IR regularization is implemented by the property
\[
\lim_{x/k^2 \to 0} R_k(x) = Z_k k^2 \quad \Leftrightarrow \quad r(y) \xrightarrow{y \to 0} \frac{1}{y}. \tag{9}
\]

The second and third properties of the regulator read
\[
\lim_{k^2/x \to 0} R_k(x) = 0 \quad \text{and} \quad \lim_{k \to \Lambda} R_k(x) \to \infty, \tag{10}
\]
and ensure that the regulator is removed in the limit \(k \to 0\) and that the initial bare action is approached at the UV scale \(\Lambda\). Moreover, the second property guarantees that the regulator vanishes for modes with \(|p| \gg k\), i.e., the theory is not affected by the regulator for large momenta.
The flow equation (5) can be mapped onto a generalized propertime representation, once the background field is identified with the full quantum field\(^1\)\cite{21,19}:

\[
\frac{\partial_t \Gamma_k[A = \bar{A}, \bar{A}]}{2 \text{STr}} = \frac{1}{2} \int_0^\infty ds \text{STr} \hat{f}(s, \eta) \exp \left( \frac{s}{k^2} \Gamma_k^{(2)} \right),
\]

where the operator \(\hat{f}(s, \eta)\) is given by\(^2\)

\[
\hat{f}(s, \eta) = \tilde{h}(s)(2 - \eta) - (\tilde{h}(s) - \tilde{g}(s))(2 - \eta) + (\tilde{H}(s) - \tilde{G}(s)) \frac{1}{s} \partial_t.
\]

In addition, we have introduced the (matrix-valued) anomalous dimension

\[
\eta := \partial_t \ln \mathcal{Z}_k = -\frac{1}{\mathcal{Z}_k} \partial_t \mathcal{Z}_k.
\]

The generalized propertime representation (11) of the flow equation has the advantage that the evaluation of the trace becomes considerably simplified.

### 2.2. Truncated Flow Equations

Even in the propertime form of Eq. (11), the flow equation cannot be solved in closed form, which necessitates further approximations. For this, we truncate the space of action functionals down to a set of operators that are considered to represent the relevant degrees of freedom for the system or at least for a particular parametric regime of the system.

For obtaining a first glance at finite-temperature gluodynamics, we concentrate on the Polyakov-loop potential, employing the simple ansatz:

\[
\Gamma_k[A, \bar{A}] = \int \frac{d^4x}{4} \left( \frac{Z_k}{4} \Gamma^{a}_{\mu\nu} F^a_{\mu\nu} + V_k \left( (v_{\mu} n^a A^a_{\mu})^2 \right) \right) + \Gamma_{gf}^k [A, \bar{A}] + \Gamma_{gh}^k [A, \bar{A}].
\]

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1. This identification itself involves an approximation for \(\Gamma^{(2)}\) in the denominator of the flow equation, as discussed in more detail in \cite{17,21,19}.

2. Terms proportional to \((h(s) - \tilde{g}(s))\) and \((\tilde{H}(s) - \tilde{G}(s))\) in \(\hat{f}(s, \eta)\) arise due to the use of the chain rule for \(\partial_t \Gamma_k\) in the flow equation (5) and manifestly represent terms arising from the spectral adjustment of the flow \(\sim \partial_t \Gamma^{(2)}\).
Here \( v_\mu \) denotes the heat-bath velocity, for which we choose \( v_\mu = \delta_{\mu 0} \). In the following, we neglect any running in the ghost- and gauge-fixing sectors, maintaining the form of Eq. (7), \( \Gamma_{gf}^{gh} = \Gamma_{A}^{gh} \). Furthermore, we neglect any other gauge-field operators except for the classical action with a wave function renormalization \( Z_k \), and, of course, the Polyakov-loop potential \( V(v_\mu A_\mu)^2 \). In order to derive the flow of the potential, it suffices to evaluate the flow for a trial background field of the simple form

\[
A_\mu^a = (n^a \phi, 0)^T,
\]

where \( n^a \) denotes a constant unit vector in color space. Furthermore, we exploit the freedom to choose suitable wave function renormalizations in the regulator function for an optimal adjustment of the regulator, cf. Eq. (8),

\[
Z_k = \begin{cases} 
Z_k^{gh} = 1, & Z_k^L = \frac{1}{\xi}, & Z_k^T \equiv Z_k 
\end{cases}
\]

for the corresponding ghost, longitudinal and transversal degrees of freedom with respect to the background field. In particular, we set the transversal wave-function renormalization equal to the background-field wave-function renormalization. The choice for \( Z_k^L \) renders the truncated flow independent of the gauge-fixing parameter \( \xi \), so that we can implicitly choose the Landau gauge \( \xi_k \equiv 0 \) which is known to be an RG fixed point [22].

It is convenient to express the flow equation in dimensionless renormalized quantities,

\[
g_k^2 = k^{d-4} Z_k^{-1} g^2, \quad \phi = \beta g \phi, \quad v_k = g^2 k^{-d} V_k,
\]

where \( \phi \in [0, 2\pi] \). Within our truncation, the flow equation in \( d = 4 \) dimensions reads

\[
\partial_t v_k(\phi) = -(4 - \eta_k) v_k(\phi) + \frac{\alpha_k}{4\pi} \sum_{n=-\infty}^{\infty} \left\{ (4 - 3\eta_k) \exp \left[ -\frac{1}{4} n^2 \left( \frac{k}{T} \right)^2 \right] \cos(n\phi) \right. \\
+ 4(2 - \eta_k) \left( \frac{T}{k} \right) \int_0^\infty dx x^2 \exp \left[ -\tilde{\omega}_n^2 - x^2 - \left( 1 - \frac{\tilde{\omega}_n^2}{\tilde{\omega}_n^2 + x^2} \right) \left( \frac{k}{T} \right)^2 \partial_\phi^2 v_k(\phi) \right] \right\},
\]

Here we have used the abbreviations

\[
\alpha_k = \frac{g_k^2}{4\pi}, \quad \eta_k = -\partial_t \ln Z_k, \quad \tilde{\omega}_n = 2\pi \frac{T}{k} n.
\]

Moreover, we have not displayed terms arising from \( \propto \partial_t \Gamma^{(2)} \), e.g. terms \( \propto \partial_t \partial_\phi^2 v_k \), on the right-hand side of the flow equation. We have dropped these terms in

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\[3\] Details of the calculation will be presented in forthcoming publication [23].
the following preliminary numerical investigation for reasons of simplicity. As a consequence, the result of Eq. (20) corresponds to a standard propertime flow \[20\]. In future work, these terms arising within the Exact RG flow will be included to facilitate a quantitative study of the differences between the Exact and the standard propertime RG in the present case. In deriving Eq. (20), we have furthermore employed a regulator which yields a particularly simple representation in propertime (or Laplace) space,

\[ \tilde{h}(s) = \delta(s - 1). \]  

At this point, it is useful to study the overlap of the present result with perturbation theory. In fact, we rediscover Weiss’s result of Eq. (4) if we (i) hold the coupling fixed, \( \alpha_k = \text{const.} \), (ii) set the anomalous dimension to zero, \( \eta_k = 0 \), and (iii) drop the complete second line of Eq. (20). The resulting simplified equation is an ordinary differential equation, that can immediately be integrated from \( k = \Lambda \) to \( k = 0 \), leading us to the perturbative one-loop result of Eq. (4).

Now, this observation helps estimating the nonperturbative content of the full Eq. (20): the occurrence of the running gauge coupling \( \alpha_k \) and the \( k \)-dependent anomalous dimension on the right-hand side signal the “RG improvement”, i.e., a resummation of an infinite set of Feynman diagrams, performed by the flow equation. Finally, the second line of Eq. (20) depends on (derivatives of) the potential itself. This term is truly nonperturbative, being induced by fluctuations of the Polyakov-loop variable on top of its own potential minimum. Since fluctuations of the degrees of freedom associated with the order parameter become highly important near the phase transition, we expect that particularly terms of this kind can cover important aspects of the nonperturbative dynamics.

In principle, the present truncation also allows for a rough calculation of the running coupling \( g_k \) including finite-temperature effects. For instance, the \( \beta \) function for the running coupling using the background field gauge reads

\[ \partial_t g_k^2 \equiv \beta_{g_k^2} = \eta_k g_k^2, \]  

and is thus related solely to the background wave function renormalization that is part of the truncation. However, important features of the running coupling in the deep infrared require much larger truncations \[13, 24\], hence we decide to take the running coupling as an external input in this work.

3. RESULTS

Let us now discuss the flow of the Polyakov-loop potential from the perturbative UV regime to the IR on the basis of our minimal approximation given by Eq. (20). In order to solve this partial differential equation, we rewrite it as an infinite set of coupled first-order differential equations by projecting it on a Fourier cosine series; this projection is naturally suggested by the fact that the \( Z(2) \) symmetry of the potential and its dependence on a compact variable allows for such a Fourier
expansion of the potential itself,

\begin{equation}
\psi_k(\varphi) = \sum_{n=0}^{\infty} \psi_k^{(n)} \cos n\varphi.
\end{equation}

With this procedure, we obtain flow equations for the dimensionless Fourier coefficients \(\psi_k^{(n)}\) by Fourier transforming the right-hand side (RHS) of Eq. (20),

\begin{equation}
\partial_t \psi_k^{(n)} = \frac{1}{\pi} \int_0^{2\pi} d\varphi \left\{ \text{RHS-Eq. (20)}[\varphi, \psi_k(\varphi)] \right\}.
\end{equation}

For a numerical evaluation of the Polyakov-loop potential, we have to truncate the infinite set at some \(n = N_{\text{max}}\) and set all \(\psi_k^{(n)}\) with \(n > N_{\text{max}}\) equal to zero. The results, shown in this work, are calculated for \(N_{\text{max}} = 2\) for simplicity. But we have confirmed that the inclusion of higher Fourier orders does not change our results significantly. As the boundary conditions for the flow equation at the UV scale \(\Lambda = 90\text{ GeV}\), we employ the result of the one-loop calculation Eq. (4).

Finally, our simple choice for the running coupling \(g_k\) in this work is

\begin{equation}
g_k \propto \left[ \ln \left( k / \Lambda_c \right) \right]^{-1},
\end{equation}

where \(\Lambda_c\) denotes a “strong scale” for which we choose \(\Lambda_c = 1\text{ GeV}\) for simplicity.

In Fig. 2, we plot the results for the flow of the coefficient \(\psi_k^{(1)}\) (left panel) and the Polyakov-loop potential at different values of \(k\) as a function of \(\varphi\) for \(T = 5 \Lambda_c\) (right panel). The flow of the first Fourier coefficient with the running coupling of Eq. (26) is compared to that for a constant coupling \(g_k \equiv \text{const.}\), and the one-loop result of Eq. (4). We find that the three different calculations are in good agreement if the

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4 The integration from \(\infty\) to \(\Lambda\) is controlled by perturbation theory, whereas the integration from \(\Lambda\) to the deep IR is controlled by the nonperturbative flow.
scale $k$ is larger than the temperature. In this regime, the potential is built up by thermal fluctuations. At $k \approx T$, the coefficient of the nonperturbative flow-equation result develops a (negative) minimum and increases to zero for even smaller $k$. Since we observe a similar behaviour also for all other Fourier coefficients, the potential flattens for $k < T$,

$$v_k^{(n)} \xrightarrow{k \to 0} 0 \quad (g_k \equiv \text{const.}) \quad \text{and} \quad v_k^{(n)} \xrightarrow{k \to \Lambda} 0 \quad (g_k \text{ of Eq. (26)}). \quad (27)$$

For the running coupling $g_k$, this behavior is depicted on the right panel of Fig. 2. By contrast, the coefficients of the one-loop result stay finite for $k \to 0$. This flatness of the potential is, in fact, nothing but the explicit manifestation of the convexity property that has to hold for the effective action in general, but is missed by perturbation theory.

Even though the observation of a convex potential is theoretically highly satisfactory, it does not tell us anything about the phase of the system. For this, the resulting expectation value of the Polyakov loop in the IR is relevant. At this point, we should stress that, even for a finally flat potential, the expectation value does not remain undetermined, but is well-defined by its $k \to 0^+$ limit (the effective potential is not flat for any nonzero $k$). However, as we can read off from Fig. 2, we do not observe a sign change of the Fourier coefficients under the flow; hence the minima of the effective potential are $Z(2)$ symmetry breaking and thus correspond to the deconfinement phase. In the present simple truncation, this holds true even for lower temperatures. Consequently, our truncation so far is only capable of describing deconfined dynamics.

4. CONCLUSIONS AND OUTLOOK

We have presented first steps towards a nonperturbative study of the Polyakov-loop potential based on an RG flow equation. Our intention so far mainly was to demonstrate the capability of our approach as a matter of principle by choosing a minimalistic approximation scheme. Already at this level, we observe important nonperturbative features, such as the backreactions of order-parameter fluctuations on top of the potential minimum and convexity of the effective potential. Confronting our first results with phenomena, we find that our truncation oversimplified the system, since we find deconfined dynamics on all scales. With hindsight, this is not too surprising, since, in addition to the Polyakov-loop dynamics, the gluon (and ghost) sector is basically approximated by its classical form, retaining “perturbative” gluon degrees of freedom all the way down to $k = 0$. From this point of view, our present minimalistic truncation leads to self-consistent results.

Future work will extend the present approach in various directions: on a technical level, the influence of the (so far neglected) terms from spectral adjustment $\propto \partial_t I^{(2)}_k$ has to be studied. This is an interesting task by its own, since it allows a quantitative analysis of the importance of these terms, distinguishing standard propertime flows from (background-field approximated) Exact RG flows. On a more conceptual level, the truncation has to be extended to include a larger
class of gluonic (and/or ghost) degrees of freedom, covering important aspects of nonperturbative dynamics and allowing for an interplay with the Polyakov-loop sector. For instance, a gluonic potential \( W_k(\frac{1}{4}F_{\mu\nu}F_{\mu\nu}) \) replacing the present simple ansatz \( \propto F_{\mu\nu}F_{\mu\nu} \) leaves room for the description of a nontrivial magnetic sector including gluon condensation as well as complex dynamics in combination with the Polyakov loop. Future work in this direction is in progress.

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