Infinite volume and continuum limits for gluon propagator in 3d SU(2) lattice gauge theory.

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(Dated: 29.04.2013)

We study the Landau gauge gluon propagator \( D(p) \) in the 3d SU(2) lattice gauge theory. We show that in the infinite-volume limit the expectation values over the Gribov region \( \Omega \) are different (in the infrared) from that calculated in the fundamental modular region \( \Gamma \). Also we show that this conclusion does not change when spacing \( a \) tends to zero.

PACS numbers: 11.15.Ha, 12.38.Gc, 12.38.Aw
Keywords: Lattice gauge theory, gluon propagator, Gribov problem, simulated annealing

I. INTRODUCTION

There are various scenarios of confinement based on infrared behavior of the gauge dependent propagators. For example, in the Gribov-Zwanziger (GZ) confinement scenario \([1, 2]\) the Landau-gauge gluon propagator \( D(p) \) at infinite volume is expected to vanish in the infrared (IR) limit \( p \to 0 \). At the same time, a refined Gribov-Zwanziger (RGZ) scenario \([3–5]\) allows a finite nonzero value of \( D(0) \). The nonperturbative lattice calculations are necessary to check the validity of each scenario as well as to check the results obtained by analytical methods, e.g., the (truncated) Dyson-Schwinger equations (DSE) approach. The DSE scaling solution predicts that the propagator tends to zero in the zero-momentum limit \([6, 7]\) in accordance with the GZ-scenario. Another - decoupling - solution \([8–11]\) allows a finite nonzero value of \( D(0) \) in conformity with RGZ-scenario.

The 3d SU(2) theory can serve as a useful test-ground to verify these predictions. It is also of interest for the studies of the high-temperature limit of the 4d theory. Last years the 3d theory has been numerically studied in a number of papers \([12–18]\). It has been shown that the propagator has a maximum at momenta about 350 ÷ 400 MeV and that zero momentum propagator \( D(0) \) does not tend to zero in the infinite-volume limit.

The Gribov copy problem still remains one of the main difficulties in computation of the gauge-dependent objects (for 3d case see, e.g., \([18]\) and references therein).

The manifold consisting of Gribov copies providing local maxima of the gauge fixing functional \( F(U) \) (defined in Section II) and a semi-positive Faddeev-Popov operator is termed the Gribov region \( \Omega \), while that of the global maxima is termed the fundamental modular region (FMR) \( \Gamma \subset \Omega \). Our gauge-fixing procedure is aimed to approach \( \Gamma \).

In paper \([20]\) it has been claimed that although there are Gribov copies inside Gribov region \( \Omega \), they have no influence on expectation-values in the thermodynamic limit i.e. for any gauge noninvariant observable \( O \)

\[
\langle O \rangle_{\Omega} = \langle O \rangle_{\Gamma}.
\]  

In our recent paper \([18]\) we attempted to check this statement. We calculated gluon propagators \( D(p) \) on different lattices (for \( p = 0 \) as well as for \( p \neq 0 \)) and then extrapolated the values of \( D(0) \) in the thermodynamic limit. It has been shown that in the thermodynamic limit \( L \to \infty \) the value of propagator \( D(0) \) clearly depends on the choice of the gauge copy. The most of our calculations in \([18]\) has been performed at \( \beta = 4.24 \) (\( a = 0.168 \) fm). The main goals of this paper are (a) to find confirmation of our observations made in \([18]\) employing different (larger) values of \( \beta \) and (b) to draw some definite conclusions about the continuum limit of the theory.

In Section II we introduce the quantities to be computed and give some details of our simulations. In Section III we present our numerical results. Conclusions are drawn in Section IV.

II. MAIN DEFINITIONS AND DETAILS OF THE SIMULATION

We consider 3d cubic lattice \( L^3 \) with spacing \( a \). To generate Monte Carlo ensembles of thermalized con-
the scale

\[ S = \beta \sum_{x,\mu > 0} \left[ 1 - \frac{1}{2} \text{Tr} \left( U_{x\mu} U_{x+1\mu} U_{x+1\mu}^\dagger U_{x\mu}^\dagger \right) \right], \tag{2} \]

where \( \beta = 4/g_B^2 a, \) \( \mu \) is a vector of unit length along the \( \mu \)th coordinate axis and \( g_B \) denotes dimensionful bare coupling. \( U_{x\mu} \in SU(2) \) are the link variables which transform under local gauge transformations \( g_x \) as follows:

\[ U_{x\mu} g(x) U_{x\mu}^\dagger = g_x U_{x\mu} g_{x+\hat{\mu}}, \quad g_x \in SU(2). \tag{3} \]

In Table I we provide the full information about the field ensembles used in this investigation. The scale is set in accordance with \[21\] where string tension is \( \sqrt{\sigma} = 440 \text{ MeV} \).

We study the gluon propagator

\[ D_{bc}^{\mu\nu}(q) = \frac{a^3}{L^3} \sum_{x,y} \exp \left( iq x + \frac{ia}{2} q (\hat{\mu} - \hat{\nu}) \right) \langle A_{b}^{\mu}(x + y + \frac{ia}{2}) A_{c}^{\nu}(y + \frac{ia}{2}) \rangle, \tag{4} \]

where the vector potentials \( A_{b}^{\mu}(x) \) are defined as follows \[22\] :

\[ A_{\mu}(x + \frac{ia}{2}) := \sum_{b=1}^{3} A_{\mu}^{b} \frac{g_{ab}}{2} = \frac{i}{g_B a} (U_{x,\mu} - U_{x,\mu}^\dagger), \tag{5} \]

and the momenta \( q_{\mu} \) take the values \( q_{\mu} = 2\pi n_{\mu}/aL, \) where \( n_{\mu} \) runs over integers in the range \(-L/2 \leq n_{\mu} < L/2\). The gluon propagator can be represented in the form

\[ D_{bc}^{\mu\nu}(q) = \begin{cases} \delta^{bc} \delta_{\mu\nu} D(0), & p = 0; \\ \delta^{bc} \left( \delta_{\mu\nu} - \frac{p_{\mu} p_{\nu}}{p^2} \right) D(p), & p \neq 0, \end{cases} \]

where \( p_{\mu} = \frac{2}{a} \sin \frac{q_{\mu} a}{2} \) and \( p^2 = \sum_{\mu=1}^{3} p_{\mu}^2 \). For \( p \neq 0 \) one arrives at

\[ D(p) = \frac{1}{6} \frac{1}{(La)^3} \sum_{\mu=1}^{3} \sum_{b=1}^{3} \langle \hat{A}_{b}^{\mu}(q) \hat{A}_{b}^{\mu}(-q) \rangle. \tag{6} \]

where

\[ \hat{A}_{b}^{\mu}(q) = a^3 \sum_{x} A_{b}^{\mu}(x + \frac{ia}{2}) \exp \left( i q x + \frac{ia}{2} \right), \tag{7} \]

and the zero-momentum propagator has the form

\[ \tilde{D}(0) = \frac{1}{9} \frac{1}{(La)^3} \sum_{\mu=1}^{3} \sum_{b=1}^{3} \langle \hat{A}_{b}^{\mu}(0) \hat{A}_{b}^{\mu}(0) \rangle. \tag{8} \]

In what follows we use the gluon propagator \( D(p) \) normalized at \( \mu = 2.5 \text{ GeV} \), so that \( p^2 D(p) = 1 \) for \( p^2 = \mu^2 \).

We employ the usual choice of the Landau gauge condition on the lattice \[22\]

\[ (\partial A)(x) = \frac{1}{a} \sum_{\mu=1}^{3} \left( A_{\mu}(x + \frac{ia}{2}) - A_{\mu}(x - \frac{ia}{2}) \right) = 0 \tag{9} \]

which is equivalent to finding a local extremum of the gauge fixing functional

\[ F_U(g) = \frac{1}{3L^3} \sum_{x,\mu} \frac{1}{2} \text{Tr} \ U_{x\mu}^g \tag{10} \]

with respect to gauge transformations \( g_x \).

To fix the gauge we choose for every gauge orbit a representative from \( \Gamma \) \[19\], i.e. the absolute maximum of the gauge fixing functional \( F(U) \). This choice is well consistent with a non-perturbative PJLZ gauge fixing approach \[23, 24\] which presumes that a unique representative of the gauge orbit needs the global extremum of the chosen gauge fixing functional. Also in the case of pure gauge \( U(1) \) theory in the Coulomb phase some of the gauge copies produce a photon propagator with a decay behavior inconsistent with the expected zero mass behavior \[23, 27\]. However, the choice of the global extremum permits to obtain the massless photon propagator.

For practical purposes, it is sufficient to approach the global maximum close enough so that the systematic errors due to nonideal gauge fixing (because of, e.g., Gribov copy effects) are of the same magnitude as statistical errors. This strategy has been checked in a number of papers on 4d and 3d theory studies for both \( SU(2) \) \[18, 28, 32\] and \( SU(3) \) \[33, 34\] gauge groups.

The gluon propagator in the deep infrared region can be reliably evaluated only when the effects of Gribov copies are properly taken into account. The gauge-fixing algorithm which we use was already successfully employed in the 4d theory at both zero \[30, 31\] and nonzero \[32, 33\] temperature. There are three main ingredients of this algorithm: powerful simulated annealing algorithm, which proved to be efficient in solving various optimization problems; the flip transformation of gauge fields, which was used to decrease both the Gribov-copy and finite-volume effects \[30, 32\]; simulation of a large number of gauge
Firstly, we extend the gauge group by the transformations (also referred to as $Z_2$ flips) defined as follows:

$$f_{\nu}(U_{x,\mu}) = \begin{cases} -U_{x,\mu} & \text{if } \mu = \nu \text{ and } x_\mu = a, \\ U_{x,\mu} & \text{otherwise} \end{cases}$$

which are the generators of the $Z_2^2$ group leaving the action (2) invariant. Such flips are equivalent to non-periodic gauge transformations. A Polyakov loop directed along the transformed links and averaged over the 2-dimensional plane changes its sign. Therefore, the flip operations combine the $2^3$ distinct gauge orbits (or Polyakov loop sectors) of strictly periodic gauge transformations into one larger gauge orbit.

We use the simulated annealing (SA), which has been found computationally more efficient than the use of the standard overrelaxation (OR) only \cite{29, 35, 36}. The SA algorithm generates gauge transformations $g(x)$ by MC iterations with a statistical weight proportional to $\exp(3V F_U[g]/T)$. The “temperature” $T$ is an auxiliary parameter which is gradually decreased in order to maximize the gauge functional $F_U[g]$. In the beginning, $T$ has to be chosen sufficiently large in order to allow traversing the configuration space of $g(x)$ fields in large steps. $T$ is decreased with equal step size. The final SA temperature is fixed such that during the consecutively applied OR algorithm the violation of the transversality condition

$$\frac{g_b a^2}{2} \max_{x,c} |(\partial A^f)(x)| < \epsilon$$

decreases in a more or less monotonous manner for the majority of gauge fixing trials until the condition (11) becomes satisfied with $\epsilon = 10^{-7}$.

To finalize the gauge fixing procedure we apply the OR algorithm with the standard Los Alamos type overrelaxation. In what follows, this method is labelled RSA (“Flipped Simulated Annealing”).

We then take the best copy (bc) out of many gauge fixed copies obtained for the given gauge field configuration, i.e., a copy with the maximal value of the lattice gauge fixing functional $F_U$ as a best estimator of the global extremum of this functional.

To demonstrate the effect of Gribov copies, we also consider the gauge obtained by a random choice of a copy within the first Gribov horizon (labelled as “fc”—first copy), i.e., we take the first copy obtained by the FSA or SA algorithms. It is instructive also to compare bc and fc propagators with the worst copy (wc) propagators which correspond to the choice of the gauge copy with minimal value of the gauge fixing functional.

### III. NUMERICAL RESULTS

To estimate the infinite-volume limit $L \to \infty$ of the zero-momentum gluon propagator $D(0; L)$ we apply the fit-formula \cite{14}

$$D(0; L) = c_1 + c_2/L.$$  \hspace{1cm} (12)

In Fig. 1 we show our values of $D(0; L)$ calculated for $\beta = 7.09$ (left panel) and $\beta = 10.21$ (right panel) for bc, fc and wc calculated by FSA method. Straight lines represent fits according to Eq. (12).

In all cases the values of $D(0)$ in the thermodynamic limit differ from zero which is in agreement with the statement made in \cite{14, 18}. Moreover, Fig. 1 demonstrates another interesting phenomenon: the Gribov copy influence remains rather strong even in the thermodynamic limit. Indeed, the infinite-volume extrapolation of $D^{bc}(0)$ differs from infinite-volume extrapolation of $D^{wc}(0)$ (as well as for $D^{wc}(0)$) which agrees with our observation made in \cite{18} for $\beta = 4.24$.

Therefore, the expectation values over the Gribov region $\Omega$ are different in the infrared from that calculated in the fundamental modular region $\Gamma$ that disagrees with the statements made in \cite{20}. This is main result of our paper.

As in \cite{18} we found that while for finite $L D^{bc}_{FS\Lambda}(0, L)$ is higher than $D^{bc}_{FSA}(0, L)$, in the infinite-volume they

| $\beta = 7.09$ ($a = 0.994$ fm) | $\beta = 10.21$ ($a = 0.063$ fm) |
|-----------------|-----------------|
| $L$ | $n_{\text{meas}}$ | $n_{\text{copy}}$ | $aL$ [fm] | $p_{\text{min}}$ [GeV] | $L$ | $n_{\text{meas}}$ | $n_{\text{copy}}$ | $aL$ [fm] | $p_{\text{min}}$ [GeV] |
| 36 | 900 | 160 | 3.38 | 0.365 | 78 | 900 | 280 | 7.33 | 0.168 |
| 48 | 900 | 160 | 4.51 | 0.274 | 92 | 1000 | 280 | 8.65 | 0.143 |
| 56 | 900 | 160 | 5.26 | 0.234 | 96 | 700 | 280 | 6.05 | 0.204 |
| 64 | 1200 | 160 | 6.02 | 0.205 |
| 78 | 900 | 280 | 7.33 | 0.168 |
| 92 | 1000 | 280 | 8.65 | 0.143 |
| 96 | 700 | 280 | 6.05 | 0.204 |

TABLE I: Values of lattice size, $L$, number of measurements $n_{\text{meas}}$ and number of gauge copies $n_{\text{copy}}$ used throughout this paper.

All details of our gauge fixing procedure can be found, e.g., in \cite{18}. For readers convenience, we will describe it shortly here.

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We find that while for finite $L D^{bc}_{FS\Lambda}(0, L)$ is higher than $D^{bc}_{FSA}(0, L)$, in the infinite-volume they...
coincide within errorbars. This remarkable agreement confirms the reliability of our estimation of \( D(0) \) in the infinite-volume limit.

For better illustration of our main result we calculated additionally the averaged difference between \( fc \) and \( bc \) propagators normalized to \( D_{bc}(p; L = \infty) \)

\[
W(p) = \frac{D_{fc}(p) - D_{bc}(p)}{D_{bc}(p; L = \infty)}
\]  

In Fig. 2 we show the dependence of \( W(0) \) on the inverse lattice size both for FSA and SA algorithms. For FSA the value of \( W(0) \) decreases for rising size, while for SA \( W(0) \) grows.

For smaller volumes the values of \( W(0) \) for SA algorithm, \( W_{SA}(0) \), are close to zero ("Gribov noise") while \( W_{FSA}(0) \) is at its maximum. The last corresponds to strong effects of flip-sectors (see, e.g., [18]). However, with increasing volume \( W_{SA}(0) \) is increasing indicating the increasing role of the copies within given flip-sector. In opposite, decreasing of \( W_{FSA}(0) \) with increasing volume implies that the role of the flip-sectors reduces.

Remarkably, in the limit \( L \to \infty \) the values \( W_{FSA}(0) \) and \( W_{SA}(0) \) coincide (within errorbars) which confirms the reliability of our fitting procedure. Results for both algorithms imply nonzero difference between \( fc \) and \( bc \) values of the propagators, this difference being \( \sim 15 \div 20\% \) in the thermodynamic limit. Thus effect is very strong.

Note that the asymptotic value \( W(0, L = \infty) \) depends weakly (if any) on the value of the spacing \( a \).

In Fig. 3 we show the momentum dependence of our \( bc \) gluon propagator \( D(p) \) for three different values of \( \beta \) (i.e., for three different values of spacing \( a \)). In all three cases the physical volumes are approximately equal with \( aL \approx 6.0 \) fm. One can see that the finite-spacing effects are very small if to compare \( \beta = 7.09 \) and \( \beta = 10.21 \) and are \( \sim 1 \div 2\% \) for \( |p| \lesssim 400 \) Mev. For larger values of \( |p| \) finite-spacing effects are even much less. We conclude that (at least) for \( \beta = 10.21 \) we can speak about the propagator in the continuum.
limit.

Note that the propagator has a maximum at nonzero value of momentum $|p| \sim 400$ MeV. Therefore, the behavior of $D(p)$ in the deep infrared region is inconsistent with a simple pole-type dependence.

In Fig. 3 we compare the momentum dependence of the $bc$ gluon propagator calculated for four different volumes for $\beta = 10.21$. Apart from $p = 0$ case, the finite-volume dependence can be seen for comparatively small values of momenta, i.e., $|p| \lesssim 0.5$ GeV. For larger values of momenta the volume dependence quickly disappears.

To compare Gribov copy effects for different values of $L$ and various momenta we define the Gribov copy sensitivity parameter $\Delta(p) \equiv \Delta(p; L)$ as a normalized difference of the $fc$ and $bc$ gluon propagators

$$\Delta(p) = \frac{D^{fc}(p) - D^{bc}(p)}{D^{bc}(p)},$$

where the numerator is the average of the differences between $fc$ and $bc$ propagators calculated for every configuration and normalized with the $bc$ (averaged) propagator.

In Fig. 4 we show the momentum dependence of $\Delta(p)$ for $\beta = 7.09$ and $\beta = 10.21$ for various volumes. As one can see, the Gribov copy influence is very strong in deep infrared and depends weakly on lattice spacing $a$. For a given value of $L$, the parameter $\Delta(p)$ decreases quickly with an increase of the momentum. It is important that for a fixed nonzero physical momentum $\Delta(p)$ tends to decrease with increasing $L$. Both observations are in agreement with the observations made earlier in [18] and for the four-dimensional $SU(2)$ theory [31].

We should emphasize that our results for $p = 0$ for all lattice spacings (see Fig. 1 and Fig. 2) as well as our results for small nonzero physical momenta presented in [18] imply that there are comparatively small nonzero momenta where Gribov copy effects also survive in the thermodynamical limit.

The last observation is essential also for the calculation of, e.g., screening masses in $4d$ theory at nonzero temperature where the momentum dependence of the gluon propagator $D(p)$ in the infrared region is important.

IV. CONCLUSIONS

We investigated numerically the Landau gauge gluon propagator $D(p)$ in the $3d$ pure gauge $SU(2)$ lattice theory. We have employed lattices with different values of $L$ for $\beta = 7.09$ ($a = 0.094$ fm) and $\beta = 10.21$ ($a = 0.063$ fm). This work is the continuation of our previous paper [18] where the most of calculations has been done for $\beta = 4.24$ ($a = 0.168$ fm).

The main goal of this work was to confirm our observations made earlier in [18] employing larger values of $\beta$ and to draw some definite conclusions about the continuum limit of the theory.

The special attention in this study has been paid to the dependence on the choice of Gribov copies. To this purpose we have generated up to 280 gauge copies for every configuration. Our $bc$ FSA method provides systematically higher values of the gauge fixing functional as compared to the $fc$ FSA and $bc$ SA methods. We stress that the choice of the efficient gauge fixing
procedure is of crucial importance in the study of the
 gluon propagator in the Landau gauge.

Our main results are the following.

1. In the thermodynamic limit \( L \to \infty \) the value of \( D(0;L) \) differs from zero. This is in agreement with RGZ-scenario and the decoupling solution of DSE, and confirms the results of numerical computations in Ref. [14, 18].

With decreasing the lattice spacing \( a \) the value of \( D(0) \) does not show the tendency to decrease.

2. The Gribov copy effects are very strong in the deep infrared region. Moreover, \( fc \) propagators do not coincide with the \( bc \) propagators even in the infinite-volume limit \( L \to \infty \) (with difference up to \( \sim 15 \div 20\% \)).

Therefore, the expectation values over the Gribov region \( \Omega \) are different in the infrared from that calculated in the fundamental modular region, i.e.,

\[
\langle O \rangle_{\Omega} \neq \langle O \rangle_{\Gamma}.
\]

3. The finite-spacing effects appear to be rather small (if to compare \( \beta = 7.09 \) and \( \beta = 10.21 \)). In the deep infrared the Gribov copy effects for \( D(p) \) depend weakly (if any) on the lattice spacing. So, we conclude that the difference between averaging over Gribov region \( \Omega \) and fundamental modular region \( \Gamma \) persists also in the continuum limit.

Acknowledgments

This investigation has been partly supported by the Heisenberg-Landau program of collaboration between the Bogoliubov Laboratory of Theoretical Physics of the Joint Institute for Nuclear Research Dubna (Russia) and German institutes, by the grant RFBR 13-02-01387. VB is supported by grant RFBR 11-02-01227-a.

References

[1] V. N. Gribov, Nucl. Phys. B139, 1 (1978).
[2] D. Zwanziger, Nucl. Phys. B364, 127 (1991).
[3] D. Dudal, S. P. Sorella, N. Vandersickel, and H. Verschelde, Phys. Rev. D77, 071501 (2008), 0711.4496.
[4] D. Dudal, J. A. Gracey, S. P. Sorella, N. Vandersickel, and H. Verschelde, Phys. Rev. D78, 065047 (2008), 0806.4348.
[5] D. Dudal, J. A. Gracey, S. P. Sorella, N. Vandersickel, and H. Verschelde, Phys. Rev. D78, 125012 (2008), 0808.0803.
[6] L. von Smekal, R. Alkofer, and A. Hauck, Phys. Rev. Lett. 79, 3591 (1997), hep-ph/9705242.
[7] R. Alkofer and L. von Smekal, Phys. Rept. 353, 281 (2001), hep-ph/0007355.
[8] J. M. Cornwall, Phys.Rev. D26, 1453 (1982).
[9] C. S. Fischer, A. Maas, and J. M. Pawlowski, Annals Phys. 324, 2408 (2009), 0810.1870.
[10] A. C. Aguilar, D. Binosi, and J. Papavassiliou, Phys. Rev. D78, 025010 (2008), 0802.1870.
[11] P. Boucaud et al., JHEP 06, 012 (2008), 0801.2721.
[12] A. Cucchieri, T. Mendes, and A. R. Taurines, Phys. Rev. D67, 091502 (2003), hep-lat/0302022.
[13] A. Cucchieri, T. Mendes, and A. R. Taurines, Phys. Rev. D71, 051902 (2005), hep-lat/0406020.
[14] A. Cucchieri and T. Mendes, PoS LAT2007, 297 (2007), 0710.0412.
[15] A. Cucchieri, A. Maas, and T. Mendes, Phys. Rev. D75, 076003 (2007), hep-lat/0702022.
[16] A. Maas, Phys. Rev. D79, 014505 (2009), 0808.3047.
[17] A. Cucchieri, D. Dudal, T. Mendes, and N. Vander-sickel (2011), 1111.2327.
[18] V. Bornyakov, V. Mitrjushkin, and R. Rogalyov, Phys.Rev. D86, 114503 (2012), 1112.4975.
[19] M. A. Semenov-tyanShanski and V. A. Franke, Zapiski Nauch. Sem. Leningradskogo old. Matematich-
eskogo inst. im. V.A.Steklova 120, 159 (1982), translation: (Plenum, NY, 1986) p.199.
[20] D. Zwanziger, Phys. Rev. D69, 016002 (2004), hep-ph/0303028.
[21] M. J. Teper, Phys. Rev. D59, 014512 (1999), hep-lat/9804008.
[22] J. E. Mandula and M. Ogilvie, Phys. Lett. B185, 127 (1987).
[23] C. Parrinello and G. Jona-Lasinio, Phys. Lett. B251, 175 (1990).
[24] D. Zwanziger, Nucl. Phys. B345, 461 (1990).
[25] A. Nakamura and M. Plewnia, Phys. Lett. B255, 274 (1991).
[26] V. G. Bornyakov, V. K. Mitrjushkin, M. Müller-Preussker, and F. Pahl, Phys. Lett. B317, 596 (1993), hep-lat/9307010.
[27] V. K. Mitrjushkin, Phys. Lett. B389, 713 (1996), hep-lat/9607069.
[28] I. L. Bogolubsky, G. Burgio, V. K. Mitrjushkin, and M. Müller-Preussker, Phys. Rev. D74, 034503 (2006), hep-lat/0511056.
[29] I. L. Bogolubsky, V. G. Bornyakov, G. Burgio, E.-M. Ilgenfritz, V. K. Mitrjushkin, and M. Müller-Preussker, Phys. Rev. D77, 014504 (2008), 0707.3611.
[30] V. G. Bornyakov, V. K. Mitrjushkin, and M. Müller-Preussker, Phys. Rev. D79, 074504 (2009), 0812.2761.
[31] V. Bornyakov, V. Mitrjushkin, and M. Muller-Preussker, Phys.Rev. D81, 054503 (2010), 0912.4475.
[32] V. Bornyakov and V. Mitrjushkin, Phys.Rev. D84, 094503 (2011), 1011.4790.
[33] V. Bornyakov and V. Mitrjushkin, Int.J.Mod.Phys. A27, 1250050 (2012), 1103.0442.
[34] R. Aouane, V. Bornyakov, E.-M. Ilgenfritz, V. Mitrjushkin, M. Muller-Preussker, et al. (2011), 1108.1735.
[35] P. Schemel, Diploma thesis, Humboldt University Berlin/Germany (2006).
[36] I. L. Bogolubsky, V. G. Bornyakov, G. Burgio, E.-M. Ilgenfritz, V. K. Mitrjushkin, M. Müller-Preussker, and P. Schemel, PoS LAT2007, 318 (2007), 0710.3234.