Dyson-Schwinger equations for IR Yang-Mills theory in different dimensions.

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Abstract. A numerical solution of the coupled Dyson-Schwinger equations for the ghost and gluon propagators in Yang-Mills theory is presented in Landau gauge. Aimed at investigating the infrared behavior of the propagators, the equations are simplified by neglecting the gluon loops, according to the ghost dominance hypothesis motivated by the Gribov-Zwanziger scenario. The equations are solved with an iterative method, eliminating the ultraviolet divergence through a continuous regulator function depending on the cut off scale. The solutions, derived for different values of the Euclidean space-time dimension, present scaling (the infrared exponents are obtained) or decoupling behavior, depending on whether the horizon condition is or not implemented. Moreover, it is shown that the running coupling constant approaches a constant value in the IR, corresponding to an attractive fixed point.

1. Introduction
In the last few decades much effort has been devoted to determine the low-energy (or long distance) properties of Green’s functions in QCD. In fact, despite of their gauge dependence, the infrared (IR) analysis of correlation functions is believed to be able to shed some light on the understanding of some mechanisms such as confinement, dynamical chiral symmetry breaking, etc.

The main difficulty of this task is the necessity to appeal to some non-perturbative technique, since the usual perturbative renormalization scheme brings to the generation of a IR Landau pole that jeopardizes the whole perturbative approach at low energies (see however [1; 2]).

Calculations on the Lattice probably represent the most reliable source of information to probe the IR sector of QCD. However, formulations on a continuum, where renormalization procedures are demanded, are employed as independent tools and in order to check the Lattice results, whose extrapolations at zero momentum present some difficulties, due to the finite nature of the Lattice.

Functional methods, in particular, like Dyson-Schwinger equations (DSEs) or Functional Renormalization Group equations (FRGs) have been extensively used to calculate (analytically or numerically) 2 or 3 point correlation functions.

These approaches have the advantage, over the “hard” numerical Lattice formulation, to better enlighten the analytical structure of the theory in the IR. However, even in the simpler pure Yang-Mills sector that is considered here, the analysis is far to be conclusive, since some truncations and approximations are required.

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A quite long debate, especially, has been pursued, on which of two qualitatively different solutions, in the Landau gauge, for the gluon and the ghost propagators are realized in the deep IR. These are known as the scaling solution, that accounts for a vanishing gluon propagator and an enhanced ghost dressing function (the ghost propagator diverges faster than its tree-level counterpart at low momentum), and the decoupling solution, where both the ghost dressing function and the gluon propagator approach finite values at vanishing momentum, with the gluon presenting therefore a massive behavior.

While the former solution is certainly more desirable, since it is compatible with the maintaining of the BRST symmetry at the quantum level and fulfills the Kugo-Ojima criterion for confinement [3; 4], it is the latter one that seems to better match the latest lattice results [5; 6]. As it has been pointed out recently [7], there is also a theoretically motivation for being more inclined towards the decoupling solution and it has to do with taking seriously into account the problem of Gribov copies [8]: the dynamics of the auxiliary fields introduced to localize the horizon function, in the Gribov-Zwanziger framework [9; 10] that constraints the configuration space of the gauge field into the first Gribov region (and that explicitly breaks the BRST symmetry), would generate condensates that affect the gluon propagator behavior, forcing it to a finite value at vanishing momentum.

Nevertheless, both solutions found for the gluon propagator account for the violation of reflection positivity [11], which is interpreted as a clue for gluon confinement.

In this contribution numerical solutions of the coupled DSEs for the gluon and the ghost 2-point functions in Landau gauge are presented in a simple truncation scheme that relies on the hypothesis of ghost dominance at low energies, that is supported by both the Gribov-Zwanziger scenario and its refined version, and is therefore supposed to give confident results at least in the deep IR region. Despite of this simplification, no particular ansatz is made regarding the functional dependence on the momentum of the propagators, since the equations are solved by iteration according to the prescription described in [12].

The solutions are calculated for different values of Euclidean space-time dimension, in order to compare them with other numerical and analytical results. Eventually, both type of solutions are found, depending on whether a boundary condition for the ghost, the horizon condition, is imposed or not. Besides, for the scaling solutions, the same IR exponents are found, in the 4 dimensional case, that were found in previous works [13; 14], with more refined truncations, and for the decoupling solution the same deep IR behavior that has been observed in the Lattice and calculated analytically [1; 2] is encountered. This suggesting that this minimal truncation scheme for the DSEs is really the only relevant for the deep IR behavior of the correlation functions.

2. Dyson-Schwinger equations in Yang-Mills theory

DSEs are a powerful tool to extract non-perturbative information about the correlation functions of a quantum field theory. They correspond to the quantum version of the classical equations of motion, since they enclose both the stationary action principle and the equal time commutation relations between the fields. In a functional formulation of quantum field theory, their derivation depends only on the assumption of the existence of a well defined integration measure, since they are based on considering as vanishing the total functional derivative inside the integral of the generating functional $Z[j]$. In a compact notation:

$$0 = \int D\varphi \frac{\delta}{\delta \varphi} e^{-S[\varphi] + j \varphi} = \left\{ -\frac{\delta S}{\delta \varphi} \left[ \frac{\delta}{\delta j} \right] + j \right\} Z[j],$$

that for the effective action $\Gamma[\phi]$, the generator of 1-PI diagrams, yields:
Taking several derivatives w.r.t. the fields, one gets exact equations for the proper vertices generated by the effective action (for a complete discussion see [11]). In particular, here the Yang-Mills action in Euclidean space-time, in Landau gauge, is considered:

\[ S = \int dx^d \left[ \frac{1}{4} F^a_{\mu\nu}(x) F^{a\mu}_{\nu}(x) + \partial_{\mu} \pi^a(x) D^{ab}_{\mu} c^b(x) + i \partial_{\mu} A^a_{\mu}(x) B^a(x) \right], \]

where the ghost term reproduce the Faddeev-Popov determinant in the gauge fixing process, and the Nakanishi-Lautrup field \( B^a(x) \) is introduced in order to constrain the vector field in the Landau gauge \( (\partial_{\mu} A^a_{\mu}(x) = 0) \).

When the DSE for the proper 2-point function of the gluon is derived, pure gluonic diagrams involving 3 and 4 gluon proper vertices appear (plus a tadpole diagram that contributes only with an infinite constant that can be absorbed in the renormalization process). In previous works these vertices have been parametrized with tensor structures, whose kernels were built by using additional physical information that could come by studying their own DSEs (that of course involve higher order vertices and must be truncated at some point) or by constraining the vertices to satisfy the Slavnov-Taylor identities (STI) that come from the BRST symmetry.

In this work such diagrams are just ignored, since it is believed, as pointed out by Gribov and Zwanziger, that the restriction of the functional integral to the region where the Fadeev-Popov operator is positive definite (the first Gribov region), which should be implemented in order to get rid of the extra gauge copies left by the Fadeev-Popov prescription, make the ghost propagator more divergent than the gluon one at zero momentum. This is also confirmed in the aforementioned refined Gribov-Zwanziger scenario [7], although the ghost dominance becomes weaker.

Moreover, from the point of view of the Wilsonian Renormalization Group (RG), as it is explained in [2], if the BRST symmetry is broken in the IR by the implementation of the horizon function, a mass term for the gluon would not be forbidden anymore, and this would change the scaling dimension of the gluon, in a way that the gluon vertices would become irrelevant operators near the IR fixed point.

It therefore seems reasonable to conjecture that these diagrams which contain only gluonic internal legs, will not give any relevant contribution to the propagators in the extreme IR. Besides, if at low energies BRST symmetry breaking has to be taken seriously, as it seems to be the case, an attempt to construct the gluon proper vertices through the analysis of the STI would not be valid anymore. More care should be taken in this way, maybe replacing the usual BRST symmetry with a more general symmetry shown by Curci-Ferrari [15] for their model that accounts for a gluon mass term.

It is noteworthy that the functional restriction to the Gribov region does not affect the form of the DSEs (that is why the horizon function was not added in (3) ), since an eventual change would come from the boundary term in (1), but this is zero due to the fact that the Fadeev-Popov operator has zero eigenvalue at the boundary (the Gribov horizon). The effect of the Gribov restriction only reveals itself in some proper boundary conditions for the ghost propagator, that at this stage are already implemented through the choice of this truncation scheme.

Finally, the truncated DSEs for the gluon and the ghost 2-point function reduce themselves to the coupled equations that are shown diagrammatically in Fig. 1. Although in the original equations referring to these diagrams, one of the ghost-gluon vertex is dressed, this has been
replaced with the bare one. Such approximation is not so drastic, since from Taylor’s argument [16] it is understood that, in Landau gauge, perturbative corrections at all orders to the proper ghost-gluon vertex are finite and scale-independent, at the symmetrical point, making the full vertex differ from the bare one only by a finite constant factor that can be absorbed in the renormalized coupling constant. Extending the approximation to the IR scales does not seem to be bold either, as it shown by the Lattice results in that region [17].

Besides, Taylor’s argument would be strictly valid only for the transversal (in the gluon momentum) part of the vertex, which is the only part that will be considered here. That is because, although the BRST breaking gives rise to a longitudinal part for the self energy of the gluon 2-point function, the gluon connected propagator continues to be transversal in Landau gauge, which is guaranteed by the DSE of the Nakanishi-Lautrup field. This means that the gluon DSE can be projected on the transversal part of the external momentum without losing any valuable information.

With this last approximation at hand, the building blocks of the diagrams are given, in momentum space, by:

$$\langle A^a_{\mu}(p)A^b_{\nu}(-q) \rangle = \frac{P^T_{\mu\nu}(p)}{\omega(p)} \delta^{ab} (2\pi)^d \delta^d(p-q), \quad \text{with} \quad P^T_{\mu\nu}(p) = \delta_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}. $$

$$\langle \epsilon^a(p)\epsilon^b(-q) \rangle = -\frac{d(p)}{p^2} \delta^{ab} (2\pi)^d \delta^d(p-q), \quad (4)$$

$$\langle 2\pi \rangle^{3d} \frac{\delta^3 \Gamma}{\delta \epsilon^a(p) \delta \epsilon^b(q) \delta A^c(k)} = P^T_{\mu\nu}(k) i g f^{abc} p_\nu (2\pi)^d \delta^d(p+q+k),$$

where a transversal projector is included in the ghost-gluon vertex, for the reason just explained. The functions $d(p)$ and $\omega(p)$ represent respectively the ghost dressing function and the inverse gluon propagator, and are the unknown functions that will be determined by numerically solving the closed system of DSEs, which are given by the following expressions:

$$\omega(p) = Z_3 p^2 + g^2 \frac{N_c}{d-1} \int \frac{d^d q}{(2\pi)^d} \frac{d(q) d(\|p+q\|)(1-(\hat{p} \cdot \hat{q})^2)}{(p+q)^2},$$

$$d^{-1}(p) = \tilde{Z}_3 - g^2 N_c \int \frac{d^d q}{(2\pi)^d} \frac{d(\|p+q\|)(1-(\hat{p} \cdot \hat{q})^2)}{\omega(q)(p+q)^2},$$

where $N_c$ is the number of colors, and the following relations have been used:
\[ f^{acd} f^{bcd} = N_c \delta^{ab}, \quad P_{\mu \nu}^T(p) = d - 1, \quad \text{(6)} \]

the second one resulting from projecting the gluon equation on the transversal part of the external momentum. \( Z_3 \) and \( \tilde{Z}_3 \) are the usual wave function renormalization constants of the gluon and the ghost respectively. Note that, according to Taylor’s argument, the vertex renormalization constant \( \tilde{Z}_1 \) has been omitted.

3. Numerical solutions

The above DSEs are solved with a cut-off regularization scheme borrowed from FRG [12]. Inside the integrals (5) the following regulator function:

\[ R(q, \Lambda) = q^2 \exp \left[ \frac{\Lambda^2}{q^2} - \frac{q^2}{\Lambda^2} \right] \quad \text{(7)} \]

is inserted in a subtracting new term, in such a way that smoothly suppresses the ultraviolet (UV) modes larger than the cut-off scale \( \Lambda \):

\[
\omega(p) = Z_3 p^2 + g^2 N_c \frac{d-1}{d-4} \int \frac{d^d q}{(2\pi)^d} (1 - (\hat{p} \cdot \hat{q})^2) \left\{ \frac{d(q) d(|p + q|)}{(p + q)^2} \right. \\
- \left. \frac{q^2}{[q^2 d^{-1}(q) + R(q, \Lambda)][(p + q)^2 d^{-1}(|p + q|)] + R(|p + q|)} \right\},
\]

\[
d^{-1}(p) = \tilde{Z}_3 - g^2 N_c \int \frac{d^d q}{(2\pi)^d} (1 - (\hat{p} \cdot \hat{q})^2) \left\{ \frac{d(|p + q|)}{\omega(q)(p + q)^2} \\
- \frac{1}{[\omega(q) + R(q, \Lambda)][(p + q)^2 d^{-1}(|p + q|)] + R(|p + q|)} \right\}.
\]

It is clear that such insertion does not affect the integrands for values of internal momentum much smaller than \( \Lambda \), because the exponential contained in the regulator function make these extra terms vanish. On the contrary, for values of momentum much larger than \( \Lambda \), the regulator function goes to zero and the extra terms become equal to the original ones, so that the integrands vanish. In this way the regulator function acts like a smooth UV cut-off.

The above system of equations is solved numerically by iteration, using as a starting guess for \( \omega(p) \) and \( d(p) \) a constant function (other trials have been made starting with other simple functions, which have given the same results) and putting the outcomes back inside the integrals, until they converge to the solutions.

Since in any dimension the integration are reduced, in spherical coordinates, to the double integrals over the modulus of the loop momentum and over the angle that this forms with the external momentum, it is immediate to perform the calculations for different values (even non integer) of dimension. The integrals are numerically evaluated using the Gauss-Legendre algorithm, up to an accuracy of about 150 nodes for each integration. The unknown functions are determined for a range of external momentum belonging to \([10^{-7} \Lambda, 10\Lambda]\).

The values of the external momentum are chosen at the zeros of Chebyshev polynomials, so that, at each step of the iteration process, inside the integrals the functions can be evaluated also for the values of momentum interpolating the Chebyshev roots, through the Chebyshev approximation, for which 120 nodes are used.

Moreover, in order to improve the convergence of the iteration process, a relaxation method is employed: at each step, inside the integrals, a linear combination of the two foregoing results
Figure 2. Scaling solutions for the gluon inverse propagator \( \omega(p) \) and the ghost dressing function \( d(p) \) at dimension \( d = 3, 3.5, 4 \).

is entered, e.g. \( \omega = r \ast \omega_n + (1 - r) \ast \omega_{n-1} \), where \( \omega_n \) is the outcome of the iteration coming before the running one, and \( r \in [0, 1] \) is a parameter that is adjusted at each step (the best choice being found at \( r \sim 0.6 \)).

It is worth to mention that an extrapolation is also needed, in the routine that evaluates the integrals, for the values of \( d(p) \) and \( \omega(p) \) at momenta that exceed the limits of the range defined for the Chebyshev approximation. In the high UV, for \( p > \Lambda \), this is done by approximating the functions with their bare counterpart (i.e. \( Z_3p^2 \) and \( \tilde{Z}_3 \)), since the integrals in the right sides of the DSEs vanish at very large momenta \( (p \gg \Lambda) \) due to the presence of the regulator function. On the other hand, in the deep IR, the values are extrapolated by doing a power-like fit (in practice a linear fit in logarithmic scale) of the functions, since it is the behavior that is expected for the propagators at a fixed point.

3.1. Scaling solution

At each step of the iteration one has to fix the renormalization constants \( Z_3 \) and \( \tilde{Z}_3 \). For the ghost, this is done, in a first moment, by imposing the horizon condition, i.e. by imposing that the ghost dressing function \( d(p) \) diverges at zero momentum. In practice this is done by subtracting the ghost DSE at zero external momentum. This is what Gribov had predicted from the effect of the restriction to the first Gribov region at the second perturbative order, and that Zwanziger confirmed at all orders.

For the gluon, the renormalization constant is fixed by imposing that in the UV region, for momenta \( p \ll \Lambda \), the gluon propagator approaches its tree-level form, \( \omega(p) = p^2 \), as a zeroth order realization of asymptotic freedom, without taking into account any resummation of perturbative corrections. This is done, at each step, by performing a quadratic fit in UV that determines the constants used to parametrize the bare gluon term:

\[
Z_3p^2 = pa^2 + m^2, 
\]

(9)

where a mass counter term is considered as generated by the gauge symmetry breaking cut-off regularization. The exact scale at which the normalization condition is implemented is not fundamental for the qualitative behavior of the solution, since (see appendix B of [14]) this truncation scheme preserves multiplicative renormalization, so that imposing the normalization condition at a different scale would only change the final result by a multiplicative factor. Moreover, in the DSEs, \( g = 1 \) and \( N_c = 3 \) are settled.
Table 1. IR exponents and running coupling constant at zero momentum.

| dimension | $\alpha$ | $\beta$ | $\alpha_R(0)$ |
|-----------|----------|---------|---------------|
| 3         | 0.60     | 0.80    | 0.76          |
| 3.5       | 0.48     | 0.99    | 1.52          |
| 4         | 0.38     | 1.19    | 2.97          |

The solutions with such renormalization scheme are depicted in Fig. 2 in a double logarithmic scale, for three different values of Euclidean dimension. As one can observe, while in the UV $d(p)$ and $\omega(p)$ present their classical behavior, coherently with asymptotic freedom, in the deep IR they exhibit a power-like dependence:

$$\omega(p)_{p \to 0} \sim p^{-\alpha}, \quad d(p)_{p \to 0} \sim p^{-\beta},$$

with the IR exponents $\alpha$ and $\beta$ that satisfy the well known sum rule (or scaling relation):

$$\alpha = 2\beta + 2 - d,$$  \hspace{1cm} (11)

that can be checked by looking at the numerical values in Table 1. The sum rule for the exponents is easily found, by simple power counting consistency, if one inserts in the truncated DSEs, with bare ghost-gluon vertices, power-like ansatze for $\omega(p)$ and $d(p)$. It is noteworthy that doing that, two different solutions for the IR exponents are found, analytically, both satisfying the sum rule [18; 19], given by the following expressions:

$$\beta = d - 2, \quad \beta \approx \frac{2d - 2}{5},$$

where the second one is an approximated formula within 2% error, and corresponds to the solutions that are found here. In fact, as pointed out in [20], this is the only solution that is compatible with both DSEs and FRGs.

Besides, by taking an $\epsilon$-expansion around the IR upper critical dimension ($d = 2$) of a Yang-Mills theory that accounts for a gluon mass term [2], one can see that the former solution of (12) corresponds to an IR repulsive fixed point, while the latter corresponds to an IR attractive fixed point (although for this last one, a fine-tuning of an irrelevant ghost operator is required, which is the equivalent of performing the horizon condition), i.e. is the only one physical relevant in the IR. Through that analysis, it is also interesting to note that such scaling solution is therefore consistent with the BRST symmetry breaking and the generation of a gluon mass term.

Taking advantage of the finite renormalization of the ghost-gluon vertex it is possible to define a dimensionless running coupling constant depending only on $\omega(p)$ and $d(p)$:

$$g_R(p) = p^{d/2 - 1} \frac{d(p)}{\omega(p)^{1/2}}.$$  \hspace{1cm} (13)

As it can be seen in Fig. 3 the running coupling constant approaches finite constant values at low momenta, for the scaling solutions, pointing out the reaching of an IR fixed point.

As before mentioned, this simple truncation scheme is thought to give confident results only in the deep IR, since in the mid-momentum region and in the UV the gluon vertices should contribute and resummed perturbative corrections should be taken into account. Nevertheless, it is comforting to recover asymptotic freedom, at least qualitatively, as it can be seen by the dumping of the running coupling at high momentum.
3.2. Decoupling solution

If the condition on the IR ghost enhancing is loosen up, i.e. the horizon condition is not explicitly implemented, a different kind of solution is numerically found, the aforementioned decoupling solution, for which the scaling relation (11) is not satisfied. The ghost wave function renormalization constant \( \tilde{Z}_3 \) is now fixed, similarly to the gluon normalization condition, by fixing the ghost dressing function to its tree-level form, i.e. \( d(p) = 1 \), at a scale near below the cut-off scale. This is achieved by performing, at each step of the iteration, a numerical fit for a constant function in this UV region.

With this renormalization scheme at hand, obviously at high momenta the same behavior as before is obtained, with the gluon propagator and the ghost dressing function approaching their bare values. However, \( \omega(p) \) and \( d(p) \) now do not diverge in the IR, but both approach finite values. It is therefore worthy to represent them only in a narrow range of deep IR momenta, to highlight in what way these functions flow towards the zero momentum point.

In Figs. 4,5 the gluon propagator \( \omega(p)^{-1} \) is depicted, in order to compare it with other numerical and analytical results. In the 3 dimensional case, a linear rise of the gluon propagator from a constant value at zero momentum can be observed. This is also seen in the lattice [5] and predicted in the \( \epsilon \)-expansion formulation near the IR fixed point [2].

This last one, ignoring the gluon kinetic term in the action, also predicts a quadratically rise in the deep IR for the gluon propagator in 4 dimension, which cannot be seen directly in Fig. 5, because it is hidden by the \( p^2 \) contribution coming from the UV fixed point. However, if the kinetic term is subtracted from the gluon propagator, the quadratically rise can be approximately observed (Fig. 6, where \( p^2 \) is put in the \( x \)-axis to make the quadratical dependence more evident).

Lastly, the running coupling constant can be seen to vanish (Fig. 7) at zero momentum, being the trivial point the only IR stable fixed point in this case [2]. As predicted in [2], \( \alpha_R \) grows from the fixed point linearly in 3 dimension and quadratically in 4 dimension.

4. Conclusions

In this work it has been argued how a minimal truncation scheme for the DSEs of the ghost and the gluon 2-point functions, where only diagrams with internal ghost legs are included, is sufficient to get the most relevant information about their behavior in the deep IR. This is justified by the IR ghost dominance effect, predicted, in Landau gauge, by the Gribov-Zwanziger scenario and its refined version.
Within such truncation scheme, where Taylor’s argument has been used to approximate the ghost-gluon vertex, two kind of solutions have been found for different dimensions: the scaling solution, where the gluon propagator and the ghost dressing function show an IR power-like behavior with exponents different from zero, and the decoupling solution, where both approach finite values at vanishing momenta.

The discriminating factor for finding one solution rather than the other resulted to be the horizon condition, i.e. imposing explicitly the divergency of the ghost dressing function at zero momentum, whether implementing or not this condition corresponds to two inequivalent renormalization schemes.

It is noteworthy that the DSEs have been solved by iteration, without any external input concerning the unknown propagators, nor without performing any angular approximation inside the loop integrals.

For the scaling solutions, the IR exponents have been evaluated and found to correspond, at least in 4 dimensions, to the same values encountered in previous works. Moreover, a running coupling constant, depending only on the 2-point functions, can be defined, and is found to approach a constant value in the IR, corresponding to a stable fixed point.

For the decoupling solution of the gluon propagator, a linear and a quadratically (if the
Figure 6. Decoupling solution for the gluon propagator in 4 dimension with the term $p^2$ subtracted.

Figure 7. Running coupling constant $\alpha_R = g_R^2/4\pi$ for the decoupling solution in 3 (left) and 4 (right) dimension.

kinetic term is subtracted) behavior has been found, in the deep IR, in 3 and 4 dimensions respectively, in agreement with the latest Lattice calculations (the quadratically behavior has been observed for the Lattice parameter $\beta = 0$) and with the $\epsilon$-expansion IR results. For these solutions, the running coupling constant vanishes at zero momentum, linearly in 3 dimensions, and quadratically in 4 dimensions.

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