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M. Lavelle
Institut für Physik
Johannes Gutenberg-Universität
Staudingerweg 7, Postfach 3980
D-6500 Mainz, Germany

*) to appear in the Proceedings of "2nd German-Chinese Symposium on Medium Energy Physics", (Bochum, Sept. 92)

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NON-PERTURBATIVE PROPAGATORS IN QCD

MARTIN LAVELLE[1]

Institut für Physik
Johannes Gutenberg-Universität
Staudingerweg 7, Postfach 3980
W-6500 Mainz, F.R. Germany

1. INTRODUCTION

Over the last two decades it has become clear that perturbation theory can only give us very limited information about QCD. For example it is not sufficient to describe that most basic of things, the mass spectrum. Although, we may hope one day to gain from the lattice approach numerical confirmation that we have the correct Lagrangian to describe hadronic physics, that day is not at hand. In the meantime it will be argued here, the operator product expansion (OPE) offers us some useful non-perturbative information about the structure of QCD.

The OPE has been extensively used in the QCD sum rules technique[SVZ76] to calculate just such fundamentals as the hadronic spectrum, coupling constants and form factors. The vacuum is here parameterised by a rather limited number of local, gauge- and Lorentz-invariant condensates, such as $< \bar{\psi}\psi >$ and $< G^2 >$, whose values are *not* directly calculated inside this formalism. For example, the value of the above quark condensate is known from PCAC and that of $< G^2 >$ was first obtained by fitting sum rules for charmonium. In general the condensates have been approximately determined by fitting to some sum rules, which ideally should be highly sensitive to the value of the condensate in question. The *universal* values of the condensates may then be employed in all other sum rules. It is noteworthy that although significant uncertainties as to the exact values of the condensates remain[Na90], QCD sum rules have given us very many predictions for hadronic quantities and this is perhaps the most successful non-perturbative method available.

More recently this approach has been extended to calculations at finite temperature and density[Ha91]. This extension is however, severely hampered by our

[1] e-mail: lavelle@vipmza.physik.uni-mainz.de

almost complete lack of knowledge as to the temperature and density dependence of the condensates, which makes the predictions of such sum rules inherently less reliable than those at zero temperature and density. In this context it is also worth remarking that there have been attempts[MR86] to parameterise the physical vacuum in a more sophisticated way by allowing the condensates a non-local distribution in space. Sum rules including such non-local vacuum expectation values are also handicapped by the need for an ansatz to describe their distribution. Reliable measurements of condensates on the lattice are unfortunately not yet available.

The sum rules are based upon the study of the two- or three-point functions of gauge-invariant currents. But the success of this method makes it natural to investigate what the OPE can tell us about those more fundamental objects the superficially divergent Greens' functions[Hae90]. Here I will discuss the OPE of the quark propagator in some detail and confine myself to a briefer treatment of the other propagators and the fundamental vertices.

Non-perturbative information on these functions would be of use in many ways. The most obvious application would be to constrain the solutions of the Schwinger-Dyson equations in QCD[Hae90]. In the lattice formulation of QCD it often proves necessary to fix a gauge in some algorithms or to measure certain operators. Knowledge of that part of the full non-perturbative structure of the two-point functions which is provided by the OPE could be of use in improving, say, the Fourier acceleration and multigrid algorithms[Ad89]. The information yielded by the OPE should also be taken into account by model builders. I shall return later to what the OPE tells us about these things. Finally a comparison of the work reported here with the field strength approach to QCD[Sch90] would also be of interest: there an interaction which seems to resemble a complicated four quark interaction at low energies and perturbative gluon exchange at high momentum transfer is found. It would be useful to investigate the appearance of power corrections as found in the OPE in this context.

The functions to be discussed here are of course gauge-dependent. This distinguishes them immediately from those considered in QCD sum rules. The analogue of gauge-invariance here is provided by its relics: BRS-invariance and the Slavnov-Tavlor identities (STI's) that follow[PT84] from it. In covariant gauges these identities imply that all corrections to the free gluon propagator must be transverse to the gluon's momentum and give us various constraints on the vertices. It is of the essence that the STI's are obeyed if gauge-invariance is not to be broken.

That earlier work[La85] on the OPE of the propagators did not fulfill the STI's provided much of the initial motivation for the work reported here. It has been assumed by many workers in this field (including initially the present author) that condensates enter the OPE of gauge-dependent functions in the same way as they enter the sum rules. This is actually not the case. Apart from where BRS-invariance

give us constraints similar to those following from gauge-invariance in sum rules, the OPE of gauge-dependent quantities will be seen to differ greatly from that of gauge-independent objects.

The structure of this talk is as follows. In Sect. 2 I discuss the choice of gauge that we must make. In particular I briefly discuss the status of non-covariant gauges and especially the Fock-Schwinger gauge. In the next section some problems with the naive application of the OPE to the QCD propagators are demonstrated. In Sect. 4 the OPE as it appears in QCD sum rules is considered in the example of the vector meson two-point sum rule. The consequences of gauge-invariance and vector current conservation are discussed. In Sect. 5 the OPE of the quark propagator is treated in some detail. It is shown that the condensates enter in a distinctly different fashion to how they appear in the sum rules. In Sect. 6 the OPE of the other QCD propagators is briefly discussed. It is shown that the STI for the gluon propagator is fulfilled and it is stressed that the ghost propagator acquires non-perturbative corrections. In Sect. 7 the OPE of the vertices is briefly reviewed. In the final section conclusions are drawn from these results. The aim of this article is to illuminate the physical phenomena evident from the OPE of the propagators. Calculational details can be found in the original papers and a compendium of results can be found in [LO92b].

2. ON CHOOSING A GAUGE

The vast majority of results reported here are from calculations made in the general class of Lorentz gauges. The gauge parameter, ξ, having the values 0 and 1 corresponding to the Landau and Feynman gauges respectively. There is I feel little need to sing the praises of this family of gauges, but it should perhaps be explained why the Fock-Schwinger (FS) gauge, $x_\mu A^\mu(x) = 0$ is not employed here, since the problems with this gauge are not widely known and it is after all a common choice in sum rule calculations. This last is because this condition leads to the following simple connection between the vector potential and the field strength:

$$A^a_\mu(x) = \int_0^1 d\alpha \alpha x^\nu F^a_{\nu\mu}(\alpha x)\,. \tag{1}$$

There are two main problems with the FS choice, perturbative and non-perturbative. The perturbative problem is concerned with the existence or otherwise of the free gluon propagator, the non-perturbative with the equivalence, for physical quantities, of this gauge to, say, the Feynman gauge.

The free gluon propagator here has been most recently considered by Modanese[Mo92]. He was unable to directly find a solution which was symmetric and fulfilled the FS gauge choice. This was in accord with earlier results[KW86]. A different, and at first sight extremely elegant, approach[Mo92], employing the above relation between the potential and the field strength, yielded a propagator which

diverges in four dimensions. This derivation however, is flawed by an incorrect handling of time ordering. The relation (1) is non-local and so the time ordering of the vector fields does not directly translate into a time ordering of the field strengths. In fact only one thing is clear: no consistent perturbative gluon propagator in FS gauge is yet known.This lack, means that the program of finding the OPE of all the superficially divergent Greens' functions cannot be carried out in this gauge.

Experience in axial gauges[CT87, LM92b] teaches us that regulating the gauge necessarily introduces other terms such as longitudinal gluons and Faddeev-Popov ghosts. Although these terms seem to vanish as the regulator, ϵ, is taken to zero, $\frac{1}{\epsilon}$ poles are generated by the loop integrals and so such terms have in general finite contributions beyond tree level. Such terms must be retained to preserve gauge-invariance. Thus it is not clear that there exists a way of obtaining a non-divergent propagator in FS gauge which would also preserve the relation (1). There are grounds to believe that non-perturbative problems with this gauge also exist. This is because Gauss' law is missing here[LM92a]. This is also missing in the temporal and light-cone gauges, and in the temporal gauge there are known cases where it is not necessary to introduce a regulator and yet the wrong answer is obtained for what should be gauge invariant quantities[La87, LM92b]. The lack of translation invariance that follows from the FS gauge condition also implies that the solutions of the SD-equations in this gauge could be rather complicated.

It thus seems clear that the FS gauge should be used very cautiously and perhaps not at all beyond tree level calculations. Since it is also not clear whether there really exists an axial gauge where ghosts and longitudinal gluons decouple[LM92b], we now return to the Lorentz class of gauges.

3. NAIVE OPE OF THE PROPAGATORS

The first work on the OPE of the propagators was by Politzer[Po76]. He calculated the coefficient of the condensate $<\bar{\psi}\psi>$ in the quark propagator in Landau gauge to leading order in α_s. This calculation was improved upon and generalised to the entire Lorentz class by Pascual and de Rafael[PR82], who obtained the self-energy correction[2]:

$$\Sigma(p) = \frac{(N_c^2-1)\pi\alpha_s <m\bar{\psi}\psi>}{2N_c^2 p^2}\Big(\frac{2\xi(2-D)}{D}\frac{\not p}{p^2} + (D-1+\xi)\frac{1}{m}\Big). \qquad (2)$$

These authors noticed the gauge dependence of the running mass term in the self-energy (2) and drew the conclusion that this could not be identified with any physical mass.

[2] Their calculation has here been further generalised to arbitrary dimension, D

Recall here that in perturbation theory the pole mass is gauge-invariant. This has been explicitly demonstrated up to two loops[Ta81]. It has been argued in a series of papers[ESS88] that this invariance persists in the OPE, when one *assumes* that the mass from the non-perturbative quark is just the pole mass and not the Lagrangian mass. Then for example at the pole, $\not{p} \approx m_{\rm Pole}$, it is clear that the gauge dependence seems to vanish in (2) in $D = 4$. Although this is formally true, it is not clear what justification there is for assuming that the mass of the soft quark is the pole mass or whether the OPE which is valid in the deep Euclidean region can be extended down to the value, $m_{\rm Pole} = 320$MeV, obtained here. Furthermore, as will be argued below, other terms are being neglected here and it must be demonstrated that they also vanish.

The OPE of the gluon propagator must obey the relevant STI and this provides a useful check on such calculations. The first work[La85] in this direction did not do this and later work[LS88] also found difficulties in fulfilling this requirement for gluonic and ghost condensates. A longitudinal part here violating the STI. Similar difficulties were later found for the three and four point gluonic vertices. It should also be noted that different results for the coefficients of $< G^2 >$ in the propagators have been given in the literature[Yn83, La85, LS88]. Clearly something is fundamentally wrong here.

Before these problems are explained, it is useful to recall how condensates enter QCD sum rules.

4. CONDENSATES AND QCD SUM RULES

Here the two point vector sum rule will be considered. This is based on the gauge-invariant correlation function

$$I_{\mu\nu}(p) = \int d^D x \quad \exp(ip\cdot x) < T(J_\mu(x) J_\nu(0)) > , \quad J_\mu(x) = \bar{\psi}(x)\gamma_\mu\psi(x) . \qquad (3)$$

Conservation of vector current implies that $I_{\mu\nu}$ must be transverse to p^μ.

Quark condensate corrections directly associated with $< m\bar{\psi}\psi >$ are easily found and are seen not to be transverse! However, one can also calculate the contribution of the condensate $< \bar{\psi} i\not{\partial}\psi >$ and one can rewrite this via the equation of motion as $< m\bar{\psi}\psi > - < g\bar{\psi}\not{A}\psi >$. This then gives the standard result, which is transverse up to the usually neglected terms of order g. The coefficient of the $< \bar{\psi} g\not{A}\psi >$ can however, also be directly calculated[LO92a]. This then gives the final result up to leading order in the coupling:

$$I_{\mu\nu} = \frac{2ig^2}{p^2}\left[\frac{4}{D} g^{\perp}_{\mu\nu}(p) < m\bar{\psi}\psi > - \left(\frac{D-4}{D} g^{\perp}_{\mu\nu}(p) + \frac{p_\mu p_\nu}{p^2}\right) < \bar{\psi}(i\not{\partial} + g\not{A} - m)\psi >\right] \qquad (4)$$

where $g^{\perp}_{\mu\nu}(p) = g_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}$. One sees that the vector current is conserved.

The point that is being made here is that in the OPE of gauge invariants only two types of condensate may enter: gauge-invariants and those that necessarily vanish. Condensates may vanish for one of two reasons, either they are equation of motion operators or they are BRS variations of other operators. In the above example and in every case to be discussed here the condensates that vanish do so as a consequence of one of the QCD equations of motion. For convenience we here give all three equations of motion:

$$\bar{\psi}(i\partial\!\!\!/ + gA\!\!\!/ - m)\psi = 0\,,$$

$$\frac{1}{2}(\partial_\mu A^a_\nu - \partial_\nu A^a_\mu)^2 + 3gf^{abc}\partial_\mu A^a_\nu A^b_\mu A^c_\nu + (gf^{abc}A^b_\mu A^c_\nu)^2 - \bar{\psi}gA\!\!\!/\psi - \bar{c}^a\Box c^a = 0\,, \quad (5)$$

$$\bar{c}^a\Box c^a - gf^{abc}\partial_\mu\bar{c}^a A^b_\mu c^c = 0\,.$$

As a consequence of this in the sum rule above, one only needs to find the coefficient of $< m\bar{\psi}\psi >$ and, say, $< \bar{\psi}i\partial\!\!\!/\psi >$. The latter may then be rewritten with the help of the equation of motion and terms of order g may be neglected in the certainty that they will cancel. This is generally done. Of course $< \bar{\psi}i\partial\!\!\!/\psi >$ must be included or vector current conservation will be violated. The neglect of such terms explains why the coefficient of $< m\bar{\psi}\psi >$ in the gluon propagator found in [La85] did not fulfill the STI.

As for gluonic condensates in the sum rule considered here, the situation is much simpler. Recall that the gluon equation of motion incorporates ghost fields with one extra factor of g. It is however, easily seen that ghosts first enter this sum rule at order g^4, thus there is no possibility for such a recombination of condensates. One merely finds the coefficient of any part of $< G^2 >$, say the abelian part $< (\partial_\mu A_\nu - \partial_\nu A_\mu)^2 >$ and from gauge-invariance one knows that this must be the coefficient of the full gluon condensate, $< G^2 >$. Of course at higher loops ghosts appear and then these effects must be taken into account.

One may now ask oneself that if all these effects depend so much upon the gauge-invariance of the sum rules or the BRS-invariance of the longitudinal part of the gluon polarisation, what does this mean for the OPE of the gauge-dependent rest of the propagators?

5. FULL OPE OF THE QUARK PROPAGATOR

In this section the coefficients of all condensates with dimension less than or equal to four that enter the quark propagator will be presented. They are given to

leading order in the coupling, so there are for example no ghost condensate contributions. They are also presented in leading order in the quark masses. The full result for the effective self-energy is then[PR82, RS86, LO92a, LO92b]

$$\Sigma(p) = \frac{\pi\alpha_s}{2N_c^2 D p^4}\Big\{(N_c^2-1)\Big(D(D-1+\xi)\frac{p^2}{m} + 2(2-D)\xi\not{p}\Big) < m\bar{\psi}\psi > +$$

$$2\Big[\Big(D^2 - 5D + 4 - N_c^2(D-1) + (2-N_c^2)\xi\Big)\not{p} + (2-D)(D-1+\xi)m\Big] < \bar{\psi}g\not{A}\psi >$$

$$-\frac{D-2}{D+2}\not{p}\Big[4N_c(D+2)p^2 < A_\mu^a A_\mu^a > -2N_c D < (\partial_\mu A_\nu^a - \partial_\nu A_\mu^a)^2 >$$

$$-4N_c(D+2) < g f^{abc}\partial_\mu A_\nu^a A_\mu^b A_\nu^c > -2(D-4) < (g A_\mu^a A_\mu^a)^2 >$$

$$+ N_c < (g d^{abc} A_\mu^b A_\mu^c)^2 > +2N_c(D+4) < (g f^{abc} A_\mu^b A_\nu^c)^2 >\Big]\Big\}. \tag{6}$$

To most easily appreciate the significance of this result, the reader should first ignore the gluonic condensates, to which we will return in a moment. The contribution of fermionic condensates to (6) should then be compared with that part given in (2). It then becomes clear that the condensates do *not* combine here to yield gauge invariant combinations and those that vanish according to equations of motion[3]. This is totally unlike the situation for the QCD sum rule considered in Sect. 4 and is a consequence of the gauge-dependence of the propagator. Clearly this behaviour should be expected to extend to all other gauge-dependent Greens' functions.

The condensates refusal to arrange themselves into $< m\bar{\psi}\psi >$, means that we cannot give accurately quantify these non-perturbative effects. Although resumming these terms may yield an effective propagator which appears to feature a running mass, we cannot give the 'mass'an accurate value. It is indeed gauge-dependent, and this dependence manifests itself not just through gauge-dependent coefficients, but also through gauge-dependent condensates. Note that it is incorrect to assume that gauge -dependent condensates have zero expectation values: this would then, for example, imply from the quark equation of motion that $< m\bar{\psi}\psi >$ vanished.

What does this tell us about the pole mass? At first sight it is hard to see how a gauge-invariant pole mass could appear from (6). However, if there really is a gauge-independent pole, then this constrains the propagator only for one value of the momentum, $p = p_{\rm Pole} = m$. This is far weaker than the equivalent constraints for all

[3] Note that the extra terms cannot be put into a gluonic equation of motion since ghost fields enter that equation of motion and there are no ghost fields in the quark propagator in this order of g^2.

external momenta in QCD sum rules and the following scenario is thinkable. Higher dimensional condensates, like $< \bar{\psi}\partial\!\!\!/ \cdots \partial\!\!\!/_n \psi >$, can be rewritten with the help of the quark equation of motion as $< m^n \bar{\psi}\psi > + O(g)$. Although these condensates are of a different dimension it could be that this summation to all orders in the mass leads to a gauge-invariant pole. This sort of resummation has been performed both for sum rules and for propagators in the past[Yn89, JM92]. For QCD sum rules this may be a useful technique, but for the propagators per se it is of dubious value. The extra terms are only superficially of higher order in the coupling, and although they will conspire to vanish in a sum rule, there is no reason to expect this to be the case in the OPE of, say, the quark propagator. Hence generally speaking, for gauge-dependent Greens' functions the only consistent expansion is one in the quark mass.

The only possible exception, is at the pole mass. Since, as discussed in Sect. 3, the validity of the OPE is open to question at such low momenta, this gauge invariance or lack of it may be only formal. It is nevertheless of interest. It should perhaps be noted that a propagator pole does not immediately imply quark deconfinement; one can construct models such that quarks in a hadron never reach their pole[LWG87].

Now consider the rôle played by the purely gluonic condensates in (6). One notices here gauge-dependent condensate combinations just as in the fermionic sector. A closer examination also reveals the appearance of the dimension two condensate[4], $< A^2 >$, and condensates which involve anticommutators of the Gell-Mann matrices! Clearly none of these can combine to yield $< G^2 >$ or vanish via the equations of motion. This further explains the difficuties mentioned at the end of Sect. 3 — assuming that the coefficients of different operators to be the coefficient of $< G^2 >$ yields various contradictory (and incorrect) results. There is also no indication from these results that gluonic condensate effects lead to any gauge invariant quark pole mass. It should also be stressed that the summing of gluonic condensate effects to all orders in the quark mass in the propagator can have no direct physical consequences.

6. FULL OPE OF THE GLUON AND GHOST PROPAGATORS

The results reported above for the quark propagator carry through to the other propagators in QCD. The gluon polarisation will therefore not be given here in full detail[LO92b]. However, the longitudinal part, which is constrained by BRS invariance to fulfill the STI, is a special case. The OPE of this is:

[4] Explicit calculation reveals that the coefficient of $< A^2 >$ in the sum rule considered in Sect. 3 is zero. From gauge invariance this must generally be the case for QCD sum rules.

$$\Pi^L_{\mu\nu} = \frac{4\pi N_c \alpha_s}{(N_c^2-1)Dp^2}\Big\{ -\frac{D(N_c^2-1)}{2N_c^2}\frac{p_\mu p_\nu}{p^2} <\bar\psi(i\partial\!\!\!/ + gA\!\!\!/ - m)\psi>$$

$$+\Big(\frac{1}{2}(2-D)\frac{p_\mu p_\nu}{p^2}\Big)\times$$

$$<\bar c^a \Box c^a - \frac{1}{2}(\partial_\mu A^a_\nu - \partial_\nu A^a_\mu)^2 - 3gf^{abc}\partial_\mu A^a_\nu A^b_\mu A^c_\nu - (gf^{abc}A^b_\mu A^c_\nu)^2 + \bar\psi g A\!\!\!/\psi>$$

$$+\frac{1}{2}D\frac{p_\mu p_\nu}{p^2} <\bar c^a \Box c^a - gf^{abc}\partial_\mu \bar c^a A^b_\mu c^c> \Big\}. \tag{7}$$

One sees that the OPE corrections indeed vanish via the equations of motion. Thus the STI is preserved. The mechanism here is just that which we saw responsible for the conservation of vector current. It is also worth noting that the ghost condensates here will naïvely be of the same order of magnitude as the gluonic ones, since they enter the gluon equation of motion.

The ghost self-energy also receives OPE corrections.

$$\Pi(p^2) = \frac{4\pi N_c \alpha_s}{(N_c^2-1)Dp^2}\Big\{\Big(1-D+\frac{1}{2}\xi(3D-5)\Big) <\bar c^a \Box c^a> + p^2 <A^a_\mu A^a_\mu>$$

$$+\frac{2-D}{2(D+2)} <(\partial_\mu A^a_\nu - \partial_\nu A^a_\mu)^2> - \frac{1}{2} <gf^{abc}\partial_\mu A^a_\nu A^b_\mu A^c_\nu> +$$

$$\frac{1}{4N_c(D+2)} <8(2(A^a_\mu A^a_\nu)^2 + (A^a_\mu A^a_\mu)^2) + N_c(2(gd^{abc}A^b_\mu A^c_\nu)^2 + (gd^{abc}A^b_\mu A^c_\mu)^2)>$$

$$-\frac{1}{2}\xi(D+1) <\bar c^a \Box c^a + gf^{abc}\partial_\mu \bar c^a A^b_\mu c^c>\Big\}. \tag{8}$$

I feel it important to stress that there is no reason to suggest, as is sometimes done, that non-perturbative corrections to ghosts may somehow be ignored. In covariant gauges at least, ghosts should not be viewed as second class citizens.

7. FULL OPE OF THE QCD VERTICES

The OPE of the vertices[LO91, AEM91, St92, Ah92] is at a less developed stage than that of the propagators. However, one moral seems to be clear: every Lorentz structure that is allowed to appear does. This may also be seen from an (incomplete) investigation of the OPE in axial gauges. Here the quark condensate contribution to the gluon polarisation was found to be[LS90]

$$\Pi_{\mu\nu}(p) = \frac{8\pi\alpha_s < m\bar{\psi}\psi >}{N_c p^2}\left(g^{\perp(1)}_{\mu\nu}(p) - g^{\perp(2)}_{\mu\nu}(p)\right),$$

$$g^{\perp(1)}_{\mu\nu}(p) := g_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}, \quad g^{\perp(2)}_{\mu\nu}(p) := \frac{g_{\mu\nu}(p\cdot\eta)^2}{p^2\eta^2} - \frac{(p_\mu\eta_\nu + p_\nu\eta_\mu)p\cdot\eta}{p^2\eta^2} + \frac{\eta_\mu\eta_\nu}{\eta^2}. \tag{9}$$

Note however, that the equation of motion has been implicitly used and that the $<\bar{\psi}g\not{A}\psi>$ terms have not been calculated. Hence the transverse part is presumably more complex than it here appears. However, the point to note is that both of the tensor structures transverse to η_μ appear. For the vertices the same phenomenon is evident, but in a more complicated form.

8. WHAT HAVE WE LEARNED?

The first conclusion must be that the condensates that enter the OPE of the propagators and vertices do so in gauge-dependent combinations. That this has not been previously recognised has been a source of confusion and error in the past.

The use of the equations of motion in the OPE is often not explicit. It is to be hoped that these results show the importance of the equations of motion. Clearly without them, the STI in Sect. 6 would not have been fulfilled.

For people working on the solutions of the SD equations, these results have a simple moral: one should use general ansätze and *not* assume that a scalar function times the perturbative propagator or vertex suffices. Additionally there is no reason to neglect non-perturbative effects in ghost Greens' functions. or for that matter to neglect ghost condensates.

I would like to conclude with a brief discussion of an application of these results. The lattice technique of Fourier acceleration introduced by the Cornell group[Ba85] to fight critical slowing down in lattice simulations of QCD requires some knowledge of the form of the two-point functions of the theory. In practice this has meant the use of a free propagator, with a numerically optimised mass. That this has not worked well for QCD, is presumably a consequence of this not being a good description of a confined propagator. The OPE results offer an alternative ansatz. The numerical uncertainty surrounding the gauge-dependent condensates in the propagators may be hoped to be there of minor significance.

Acknowledgements: I would like to thank my collaborators as well as the organisers and participants of this stimulating conference. I also thank K. Schilcher for a discussion and the Graduierten Kolleg of the University of Mainz for financial support.

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