On the uniqueness of the Standard Model of particle physics

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Abstract

The running strong coupling $\alpha(Q)$, and the gluon propagator from QCD have been compared to similar quantities in the quanton model, a generalisation of QED with massless fermions (quantons) and scalar coupling of boson fields. In the latter model a series of bound states (which can be related to different flavours) have been obtained. Assuming a weighting of their momentum distributions with the average momentum $\tilde{Q}_i$ of each state, a running of the coupling $\alpha(Q)$ is obtained, which is in quantitative agreement with $\alpha_s(Q)$ from QCD. Also with a similar weighting the gluon propagator from lattice QCD simulations is well described. This indicates clearly that QCD and thus the Standard Model is not a unique description of fundamental interactions.

Different from the Standard Model, the quanton model is simple and yields bound states with correct masses. This may indicate that this model yields a more realistic description of fundamental forces.

PACS/keywords: 11.15.-q, 12.40.-y, 14.40.-n/ Comparison of the running coupling $\alpha(Q)$ and the gluon propagator of QCD with the quanton model, an extension of QED with scalar coupling of gauge bosons. Equivalent description in both models indicating that the Standard Model is not unique.

From the ultimate theory of fundamental forces one expects a structure, which is simple, complete and unique. Whether nature reduces the fundamental forces to a very simple form is not known and has to be found out, but completeness can be tested. Apart from the fact that the theory has to describe the wealth of existing data, it is important that it couples to the absolute vacuum with energy $E_{\text{vac}} = 0$. The uniqueness can be tested only by comparing the theory with other models.

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The Standard Model of particle physics (SM) is composed of different gauge theories, quantum electrodynamics (QED), weak interaction theory, and quantum chromodynamics (QCD), and has been applied successfully to the description of hadronic and leptonic processes over the last decades. It is certainly not simple and also not complete, since the elementary fermions (quarks and leptons) are massive. Another shortcoming of the SM is that in the relativistic theory stationary states of elementary fermions, the most important features of particle physics, cannot be described. This has been taken less seriously, since bound states can be described in non-relativistic approximations. However, this is not possible for the strong interaction, which has a large coupling strength.

To understand the mass problem of elementary fermions in the SM, the Higgs-mechanism has been adopted, in which the g.s. of elementary particles is lowered by symmetry breaking. However, this mechanism implies the existence of heavy Higgs-bosons, which have never been found in spite of extensive searches over 20 years. Another severe problem of the Higgs-mechanism is that it requires an enormously high energy density of the vacuum in striking disagreement with astrophysical observations.

Recently, a new formalism (quanton model) has been developed for the description of fundamental forces, based on a different understanding of mass, given solely by binding effects of elementary massless fermions (quantons). This model can be used to test the uniqueness of the SM. In the present paper two basic quantities of QCD - which have been investigated extensively - have been compared with corresponding properties of the quanton model: the running strong coupling \( \alpha_s(Q) \) and the ‘gluon propagator’, the latter taken from lattice gauge simulations.

Starting from QED with scalar coupling of two boson fields \( A_\mu A^\mu \), the Lagrangian may be written in the form

\[
\mathcal{L} = \frac{1}{m^2} \bar{\Psi} i\gamma_\mu D_\mu (D_\nu D^\nu) \Psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu},
\]

where \( m \) is a mass parameter, \( \Psi \) a massless fermion (quanton, q) field, \( D_\mu = \partial_\mu - igA_\mu \) the covariant derivative, and \( F_{\mu\nu} = \partial^\alpha A_\nu - \partial^\nu A_\mu \) the field strength tensor. In this Lagrangian higher derivatives of the fermion field appear, which have to be removed. This is possible by a symmetry between fermions and one boson field \( \partial^\nu - igA_\nu = const \), leading to a constrained Lagrangian of the form

\[
\mathcal{L} = \frac{1}{\bar{m}} \bar{\Psi} i\gamma_\mu D_\mu D_\nu \Psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu},
\]  

\( (1) \)

\(^2\)with insignificant mass with respect to the mass scale of particle physics
The first term of this Lagrangian gives rise to the following two terms

\[ \mathcal{L}_{2g} = -\frac{ig^2}{\tilde{m}^2} \bar{\Psi} \gamma_\mu A^\mu (\partial_\nu A^\nu) \Psi \]  

(2)

and

\[ \mathcal{L}_{3g} = -\frac{g^3}{\tilde{m}^2} \bar{\Psi} \gamma_\mu A^\mu (A_\nu A^\nu) \Psi . \]  

(3)

These two parts of the Lagrangian do not have derivatives of the fermion field. Further, the derivative of \( A^\nu \) appears in \( \mathcal{L}_{2g} \), but this field appears in \( \mathcal{L}_{3g} \) only in combination with another boson field. Therefore, the problems of higher derivative Lagrangians are avoided. Generally, a Lagrangian with a mass denominator \( \tilde{m}^n \) (with \( n \geq 1 \)) leads to a non-renormalizable theory, however, in the present case the Lagrangian yields finite matrix elements and consequently results in full space.

From these Lagrangians ground state matrix elements \( \langle g.s. \mid K(\tilde{p} - \tilde{p}) \mid g.s. \rangle \) have been derived, where \( K(\Delta \tilde{p}) \) are the multi-boson field operators of the above Lagrangians (2) and (3). These yield contributions only, if the two boson fields overlap in space and time.

By equal time requirement of the overlapping boson fields a reduction to three dimensions is possible, which gives rise to two potentials, which are given in r-space by

\[ V_{2g}(r) = \frac{\alpha^2 \hbar^2 \tilde{E}^2}{2m^3} \left( \frac{d^2 w(r)}{dr^2} + \frac{2}{r} \frac{dw(r)}{dr} \right) \frac{1}{w(r)} , \]  

(4)

and

\[ V_{3g}(r) = \frac{\hbar}{m} \int dr' \rho(r') V_g(r - r') , \]  

(5)

where \( w(r) \) and \( \rho(r) = w^2(r) \) are wave function and (quasi) density\(^3\) of a two-boson field, \( \tilde{E} \) the mean energy of scalar excitation in the potential (4), and \( V_g(r) \) a boson-exchange potential given by \( V_g(r) = -\alpha^3 \hbar \frac{f(r)}{r} \).

The potential \( V_{2g}(r) \) corresponds to the ‘confinement’ potential required in hadron potential models \([5]\), whereas \( V_{3g}(r) \) is related to the usual boson-exchange potential derived from basic gauge theories, but due to its more complex structure it is finite for \( r \to 0 \) and \( \infty \) and scales with the coupling strength \( \alpha^3 \).

For a \( q\bar{q} \) system in a scalar state \( (J^\pi = 0^+) \), angular momentum \( L=1 \) is needed. Therefore, for this case a p-wave density is required in eq. (5), which is related to \( \rho(r) \) by

\[ \rho^p(\vec{r}) = \rho^p(r) Y_{1,m}(\theta, \Phi) = (1 + \beta R \frac{d}{dr} \rho(r)) Y_{1,m}(\theta, \Phi) . \]  

(6)

\(^3\rho(r) \) has dimension \( fm^{-2} \) due to the dimension \( GeV^2 \) of the two-boson field.
\( \beta R \) is determined from the condition \( \langle r_{\rho^s} \rangle = \int d\tau \ r \rho^s(r) = 0 \) (elimination of spurious motion). This yields a boson-exchange potential given by

\[
V_{3g}^s(r) = \frac{\hbar}{m} \int d\vec{r}' \rho^s(\vec{r}') \ Y_{1,m}(\theta', \Phi') \ V_g(\vec{r} - \vec{r}') = 4\pi \frac{\hbar}{m} \int dr' \rho^s(r') \ V_g(r - r') .
\] (7)

By requiring that the boson-exchange force can act only within the two-boson density \( \rho(r) \), which gives rise to the constraint

\[
V_{3g}^s(r) = c_{\text{pot}} \rho(r) ,
\] (8)

the density as well as the parameters of the interaction cut-off function \( f(r) = (e^{(ar)^\kappa} - 1)/(e^{(ar)^\kappa} + 1) \) \( e^{-cr} \) can be deduced self-consistently, see the details in ref. [3]. This yields

\[
\rho(r) = \rho_o \ \left[ \exp\{-(r/b)^\kappa\}\right]^2 \ \text{with} \ \ \kappa \sim 1.3 - 1.5 .
\] (9)

Inserting this form of \( \rho(r) \) in \( V_{2g}(q) \) leads to the explicit form

\[
V_{2g}(r) = \frac{\alpha^2\hbar^2 \tilde{E}_2^\kappa}{m^3} \left[ \frac{\kappa}{b^2} \left( \frac{r}{b} \right)^{\kappa - 2} \left[ \kappa \left( \frac{r}{b} \right)^\kappa - (\kappa + 1) \right] \right] .
\] (10)

The mass of the system is defined by

\[
M_n = -E_{3g} + E_{2g}^n ,
\] (11)

where \( E_{3g} \) and \( E_{2g}^n \) are the binding energies in \( V_{3g}^s(r) \) or \( V_{3g}^v(r) \) and \( V_{2g}(r) \), respectively, calculated by using a mass parameter \( \tilde{m} = 1/2 \tilde{M} \), where \( \tilde{M} \) is the average mass of the system, weighted over vector and scalar states. However, since this weighting is not known, \( \tilde{m} = 1/2 \tilde{M} \) is used, where \( \tilde{M} \) is the ground state mass of the system. This allows to use the additional constraint \( M = M_1 \). In this way, the mass contributions due to excited states is included in \( \tilde{E} \), which is used as fit parameter. The coupling constant \( \alpha \) is obtained by matching the mass to the lowest binding energy in eq. (11).

The binding energies in \( V_{3g}^s(r) \) are negative. Using the energy-momentum relation in the form \( E_{\text{vac}} = 0 = \sqrt{<Q^2_\rho>} + \tilde{E}_{3g} \), where \( <Q^2_\rho> \) is the mean square momentum of \( \rho(r) \) and \( \tilde{E}_{3g} \) a weighted average of \( E_{3g} \) between vector and scalar states, this yields another constraint

\[
\tilde{E}_{3g} = -\sqrt{<Q^2_\rho>} .
\] (12)

Different from the binding in \( V_{3g}^s(r) \), which does not correspond to real mass generation, the binding energy \( E_{2g} \) is positive and allows creation of stable particles out of the absolute
vacuum of fluctuating boson fields, if two rapidly fluctuating boson fields overlap and trigger a quantum fluctuation with energy $E_{2g}$.

Using the constraints (8) and (12) and the energy-mass relation (11), all parameters of the model are determined (with some ambiguities) for a given slope parameter $b$. This last parameter may be determined by the structure of the vacuum. For mesonic systems this is discussed in ref. [10] and leads to a description, in which all parameters of the model are fixed. This is based on a vacuum potential sum rule, assuming a global boson-exchange interaction in the vacuum $V_{\text{vac}}(r) \sim 1/r^2$. Further, the different potentials $V_{3g}^i(r)$ (where $i$ are the discrete solutions) sum up to $V_{\text{vac}}(r)$

$$
\sum_i V_{3g}^i(r) = V_{\text{vac}}(r) = \tilde{f}_{as}(r)(-\tilde{\alpha}_3^3 \hbar r_o/r^2) e^{-\tilde{\epsilon}r},
$$

where $\tilde{f}_{as}(r)$ and $e^{-\tilde{\epsilon}r}$ are cut-off functions similar to those for the boson-exchange interaction discussed above.

A sum rule analysis similar to that discussed in ref. [10] has been performed, but for the present study the higher energy region up to several hundred GeV has been included. A comparison of the resulting boson-exchange potentials with the sum rule (13) is given in fig. 1 with the deduced masses and parameters in table 1. In addition to the states discussed in ref. [10] another solution is obtained in the hundred GeV region, with a vector state at about 91 GeV and a scalar state at the mass of the observed $t\bar{t}$ system; therefore, these two states can be identified with a 'top' flavour system. Interestingly, the vector state is at the mass of the $Z^0$ boson, consequently, this particle has to be identified in our model as part of the 'top' system.

To make a comparison with the running coupling strength $\alpha^{QCD}(Q)$ one has to go to the Fourier transform of the vacuum potential $V_{\text{vac}}(r)$ (13), which can be written in the form $V_{\text{vac}}(Q) = \alpha^3(Q)/Q = \sum_i V_{3g}^i(Q)$. This leads to $\alpha(Q) = [\sum_i V_{3g}^i(Q)]^{1/3} Q^{1/3}$. Individual coupling functions $\alpha_i(Q)$ can also be defined by $\alpha_i(Q) = [V_{3g}^i(Q)]^{1/3} Q^{1/3}$. For the comparison with QCD the following form is used

$$
\alpha^{QCD}(Q) = [\sum_i w_i V_{3g}^i(Q)]^{1/3} Q^{1/3}
$$

with weighting factors $w_i$, which are adjusted to the QCD data in fig. 2. A quantitative agreement with $\alpha^{QCD}(Q)$ is obtained in the momentum range covered by the solutions in
Table 1: Deduced masses (in GeV) of scalar and vector $q^+q^-$ states in comparison with the lowest $0^{++}$ and $1^{--}$ mesons \[1\].

| Solution | (meson) | $M$   | $M^{exp}$ |
|----------|---------|-------|----------|
| 1        | scalar  | $\sigma$ | 0.55 | 0.60±0.2 |
| 2        | scalar  | $f_0$  | 1.70 | 1.70±0.2 |
|          | vector  | $\omega$ | 0.78 | 0.78     |
| 3        | scalar  | $f_0$  | 3.28 |          |
|          | vector  | $\Phi$ | 1.02 | 1.68±0.02 |
| 4        | scalar  | not seen | 12.7 |          |
|          | vector  | $J/\Psi$ | 3.10 | 3.097    |
| 5        | scalar  | not seen | 40.4 |          |
|          | vector  | $\Upsilon$ | 9.46 | 9.46     |
| 6        | scalar  | top    | \(\sim 370\) | \(\sim 370\) |
|          | vector  | $\,(Z^0)\,$ | \(\sim 91\) | 91.2     |

| Sol.| $\kappa$ | $b$   | $\alpha_e$ | $c$   | $a$   | $\langle r_\rho^2 \rangle$ | $\langle Q_\rho^2 \rangle^{1/2}$ | $\tilde{E}$ |
|-----|---------|-------|-------------|-------|-------|----------------|----------------|---------|
| 1   | 1.50    | 0.831 | 0.257       | 2.24  | 5.8   | 0.761          | 0.59            | 1.0     |
| 2   | 1.46    | 0.264 | 0.260       | 7.20  | 18    | 0.080          | 1.57            | 0.6     |
| 3   | 1.44    | 0.149 | 0.281       | 13.2  | 30    | 0.026          | 2.98            | 0.5     |
| 4   | 1.40    | 0.054 | 0.327       | 33    | 82    | 0.36 $10^{-2}$ | 7.73            | 0.8     |
| 5   | 1.36    | 0.020 | 0.390       | 95    | 220   | 0.51 $10^{-3}$ | 20.7            | 1.5     |
| 6   | 1.30    | 0.0047| 0.714       | 390   | 900   | 0.32 $10^{-4}$ | 90.8            | 4.4     |
table 1 assuming $w_i = \frac{1}{3} < Q_i^2 >^{-1/2}$. This is shown in fig. 4, which displays eq. (14) by solid line as well as the individual components by dashed and dot-dashed lines.

From this comparison one can draw the following conclusions: first, the correspondence with QCD supports the concept of a vacuum potential sum rule in the present model with a sequence of solutions consistent with different flavour states. It is evident that $\alpha(Q)$ is nothing else but a different form of the vacuum potential sum rule with a strong weighting of the high momentum region. One can see that the 'top' system at about 100 GeV is needed. Further, the abrupt fall-off beyond $Q=100$ GeV in our calculations indicates that the flavour states are not limited to the states presently known but continue to larger masses, with the next flavour state in the one TeV mass region. Second, the factor $1/3$ in $w_i$ is probably due to colour, whereas the factor $\tilde{Q} = < Q_i^2 >^{-1/2}$, which leads to the strong fall-off of $\alpha^{QCD}(Q)$, appears to be due to the non-Abelian character of QCD. If one relates $\alpha(Q)$ to the mass, the quanton model is correct, whereas $\alpha^{QCD}(Q)$ suffers from increasing mass deficits for larger Q-values, apparently related to the strong increase in the needed quark masses.

Another problem of QCD is that only perturbative solutions exist, restricting QCD analyses to reactions with large momentum transfers, as deep inelastic scattering. For this type of reaction, the extracted strength has to be weighted with the probability of the momentum transfer $Q$, which is $1/Q$. Therefore, the quanton model yields the same result as QCD, if for deep inelastic reactions the strength is weighted with $1/Q$.

For a second comparison with QCD at much smaller momenta, one has to go to lattice gauge theory, which is not limited to perturbative solutions. In this approach hadron masses have been calculated quite successfully, but also a good description of the confinement potential [12] has been obtained. In particular, the 'gluon propagator' has been investigated quite intensely, see e.g. refs. [6] [7] [8] [9], since its structure lacks a simple direct understanding. This quantity is not gauge independent, a summary of lattice calculations in different gauges [7] [8] [9] is shown in fig. 5. In our framework we describe the corresponding 'boson propagator' $P_g$ as a sum of boson-exchange potentials over $i$ flavour states [10] $P_g = \sum_i V_{3g}^i(Q)$. Similar to the analysis of $\alpha^{QCD}(Q)$ we replace $P_g$ by

$$P_g^{QCD} = \sum_i w'_i V_{3g}^i(Q)$$

(15)

with weighting factors $w_i'$ fitted to the lattice data. Good fits are shown in fig. 5 with
average weightings $w'_i \sim 33 \cdot \frac{1}{2} < Q_i^2 >^{-1/2}$ and individual contributions given by the dot-dashed lines. So, we observe a similar $< Q_i^2 >^{-1/2}$ dependence as found for $\alpha^{QCD}(Q)$.

Concerning the weak interaction part of the SM, this interaction may be replaced in the quanton model by the spin-spin interaction between quantons, assuming a structure of leptons as bound states of three quantons, see ref. [11]. It is interesting to note that the spin-spin force for neutral elementary particles is not considered in the SM. This interaction has properties very similar to the weak interaction, a coupling strength many orders of magnitude smaller than the electromagnetic force and consequently an extremely short range [11]. However, the crucial difference to the weak interaction theory is that the spin-spin force does not require massive gauge bosons and (due to the Higgs-mechanism) a large energy density of the vacuum.

In conclusion, calculations of the momentum dependence of the coupling $\alpha(Q)$ and the boson propagator within the quanton model have been compared with the corresponding quantities from QCD. A good agreement of these quantities is obtained in both theories, if the results of the quanton model are scaled with the momentum. This indicates clearly that QCD and thus the SM is not the unique theory of the strong interaction.

On the other hand, the quanton model appears to have all necessary properties of a fundamental theory, it has a very simple structure with only two quantons, charged and uncharged, coupled by the same gauge boson to the charge and spin, respectively. Further, it may to be complete, but further test are needed before this model can be regarded as a realistic description of fundamental forces.

We thank B. Loiseau for his continuous support during the development of the model.

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Figure 1: Boson-exchange potentials for different solutions (given by dot-dashed and dashed lines) and sum given by solid line. This is compared to the vacuum sum rule \([13]\) given by the dot-dashed line overlapping the solid line. A pure potential \(V = -\text{const}/r^2\) is shown also by the lower dot-dashed line. The lower part shows the same lines but with a vertical scale enlarged to -140 GeV.
Figure 2: Momentum dependence of the coupling strength $\alpha(Q)$ from QCD analyses (solid triangles) and lattice QCD simulations (open triangles) in comparison with our results. Applying for all solutions a normalisation $\frac{1}{3} < Q_i^2 >^{-1/2}$, the dot-dashed and dashed lines correspond to the individual solutions [10] with an additional solution in the 100 GeV region. The sum is given by the solid line, which is in good agreement with QCD.
Figure 3: Data on the gluon propagator from lattice gauge calculations in different gauges \[7, 8, 9\] in comparison with calculations within our model using a sum of the different flavour contributions given in ref. \[10\], which yields a good description of the lattice data (solid lines). Individual components with average weighting \(\sim 33 < Q_i^2 >^{-1/2}\) are shown in the upper and lower part by dot-dashed and dashed lines.