6. Field theories of matter and gauge interactions

The most characteristic fact of particle physics is that the interactions among fermionic matter building blocks is mediated by gauge bosons (e.g. the photon) (NB: the Higgs boson is somewhat Janus-faced, it carries matter properties as well as mediates a force via Yukawa interactions). The underlying local gauge symmetry that we have already encountered in Maxwell's theory is largely responsible for the resulting structures. In G. 't Hooft's words: we are "under the spell of the gauge principle".

6.1 (Quantum) electrodynamics (QED)

Starting from the Maxwell Lagrangian (1.60) known from classical electrodynamics,

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - J_\mu A^\mu , \qquad (6.1)$$

let us try to add fermionic electron/positron degrees of freedom in the form of a Dirac spinor field $\psi(x)$, while preserving the local gauge

Symmetry under gauge transformations:

$$A_\mu(x) \rightarrow A_\mu(x) + \partial_\mu \Lambda(x) \quad , \Lambda(x): \text{arbitrary} \qquad (6.2)$$

Assuming that the interaction can be written in terms of a suitable choice for the source $J_\mu = J_\mu[\bar\psi, \psi]$, the action remains invariant, if

$$S = \int d^4x \left[-\frac{1}{4} \underbrace{F_{\mu\nu}F^{\mu\nu}}_{\substack{\text{gauge}\\ \text{invariant}}} - J_\mu A^\mu \right]$$

$$\rightarrow \int d^4x \left[-\frac{1}{4} F_{\mu\nu}F^{\mu\nu} - J_\mu A^\mu - J_\mu \partial^\mu \Lambda \right]$$

$$\overset{\text{i.b.p}}{=} \int d^4x \left[-\frac{1}{4} F_{\mu\nu}F^{\mu\nu} - J_\mu A^\mu + \Lambda \partial^\mu J_\mu \right] \qquad (6.3)$$

the source is conserved, $\partial^\mu J_\mu = 0$.

Indeed, the free Dirac theory

$$S_D = \int d^4x \left[\bar\psi i\partial\!\!\!/ \psi - m\bar\psi\psi \right] \qquad (6.4)$$

offers a conserved source: the Noether current j^μ associated with $U_V(1)$ vector symmetry

$$\psi \rightarrow e^{i\alpha}\psi \quad , \bar\psi \rightarrow \bar\psi e^{-i\alpha} \qquad (6.5)$$

We have determined the resulting Noether current in the exercises:

$$j^\mu = \bar{\Psi} \gamma^\mu \Psi \qquad , \partial_\mu j^\mu = 0 \tag{6.6}$$

This suggest to identify J^μ with the Noether current,

$$J^\mu = e\, j^\mu , \tag{6.7}$$

where we have allowed for a coupling constant e that parametrizes the strength of the interactions between the Maxwell and the Dirac field.

Upon insertion of (6.7) into (6.1) and adding the Dirac action (6.4), we arrive at

$$S_{QED} = \int d^4x \left[-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \bar{\Psi}\, i \not{D}[A] \Psi - m \bar{\Psi} \Psi \right] \tag{6.8}$$

where we have used the covariant derivative (c.f. (3.34))

$$D_\mu[A] = \partial_\mu + i e A_\mu \qquad \text{and } \not{D} = \gamma^\mu D_\mu . \tag{6.9}$$

Eq. (6.8) denotes the classical action of Quantum Electrodynamics (which becomes "Quantum", of course, only upon quantization of the fields).

Our construction guarantees, that S_{QED} is invariant under the local gauge symmetry (6.2) as well as the global vector symmetry (6.5) separately.

However, the interesting observation now is that S_{QED} is fully invariant under a simultaneous local transformation of both fields:

$$\begin{aligned} A_\mu(x) &\to A_\mu(x) + \partial_\mu \Lambda(x) \\ \psi(x) &\to e^{-ie\Lambda(x)}\,\psi(x) \\ \bar\psi(x) &\to \bar\psi(x)\, e^{ie\Lambda(x)} \end{aligned} \tag{6.10}$$

This is the same typo of local $U(1)$ symmetry that we have already encountered for the abelian Higgs model in (3.33) and ff.

The essential building block is the covariant derivative $D_\mu[A]$, which guarantees that

$$\left(D_\mu[A]\,\psi\right) \to e^{-ie\Lambda(x)}\left(D_\mu[A]\psi\right) \tag{6.11}$$

– despite the partial derivative – transforms with a simple $U(1)$ phase factor.

Already on this "classical" level, the theory (6.8) is useful, as (together with a proton field) it offers the relativistic version of the quantum mechanical hydrogen-problem, describing relativistic effects in atomic physics rather accurately (cf. your course on advanced quantum mechanics).

QED, however, celebrates its greatest successes in the quantized version for the quantitative description of the anomalous moment of the electron or Lamb shift effects in atoms.

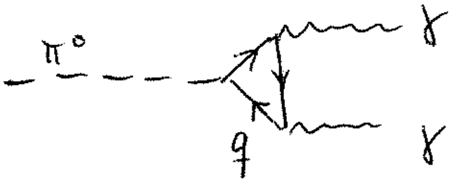
Here, we plan to go beyond and wish to view this theory as a first simple example of a gauge theory.

6.2 (Quantum) Chromodynamics

The necessity of a further quantum number, i.e., another type of charge for elementary constituents became clear from the observation of Baryon resonances with three quarks in the same flavor and spin state

$$\begin{aligned} |\Delta^{++}\rangle &= |u\uparrow\rangle\,|u\uparrow\rangle\,|u\uparrow\rangle \\ |\Omega^{-}\rangle &= |s\uparrow\rangle\,|s\uparrow\rangle\,|s\uparrow\rangle \\ |\Delta^{-}\rangle &= |d\uparrow\rangle\,|d\uparrow\rangle\,|s\uparrow\rangle \end{aligned} \tag{6.12}$$

seemingly contradicting Pauli's exclusion principle. Upon adding a further quantum number, the required antisymmetrization for the fermionic constituents can be realized with respect to this new quantum number, called "color". As a consequence, processes which can proceed via different internal values of this quantum number become proportional to it, e.g. pion decay into two photons,



$$\sim \quad N_c \tag{6.13}$$

According to QFT, the decay proceeds via an internal quark fluctuation. As the quarks now can occur in diffently colored versions, the process is proportional to the "number of colors" N_c. The experimental result is $N_c = 3$. i.e. in addition to the different quark "flavors" $f = u, d, s, c, b, t$ quarks also carry a "color" index $i = 1, 2, 3$:

$$\psi(x) = \psi_f^i(x) \tag{6.14}$$

In the following, we ignore the flavor and concentrate on the color index $i = 1, 2, 3$.

The above experiments suggest that there is at least a global symmetry in an internal color space by which we can transform the spinors:

$$\psi^i \Rightarrow \psi'^i = U^{ij}\psi^j \quad . \qquad (6.15)$$

The decisive aspect of this symmetry exerting a strong influence on the resulting dynamics, however, is that this symmetry turned out to be a <u>local</u> symmetry analogous to the one of QED:

$$\psi'^i_{(x)} = U^{ij}(x)\, \psi^j_{(x)} \qquad (6.16)$$

where $U(x) \in SU(N_c)$ is a matrix, being an element of the matrix group $SU(N_c)$, i.e. complex unitary matrices with $\det U = 1$.

This local symmetry property cannot be read off from kinematical observations as the ones given above, but require a close look at the dynamics or bound-state spectra of the system.

Let us first recall a few basic facts about the Lie groups $SU(N_c)$ and their corresponding Lie algebra. The complex $N_c \times N_c$ matrices U^{ij} with

$$U^\dagger U = \mathbb{1} = U U^\dagger \quad , \det U = 1 \tag{6.17}$$

form a representation of $SU(N_c)$. The exponential map

$$U = e^{iH} \quad \text{where } H = H^\dagger \text{ hermitean } N_c \times N_c \text{ matrix} \tag{6.18}$$

parametrizes U in terms of $\underbrace{N_c^2 - 1}_{\det U = 1}$ real parameters. (6.19)

This implies that H can be spanned by $N_c^2 - 1$ linearly independent hermitean matrices which serve as generators of $SU(N_c)$:

$$U = e^{-i\omega_a T^a} \quad , (T^a)^{ij} : \text{ generators of } SU(N_c) \quad i,j = 1,\dots,N_c \quad a = 1,\dots,N_c^2-1 \tag{6.20}$$

where ω_a are real parameters, and the T^a can be chosen tracefree since

$$1 = \det U = \det e^{-i\omega_a T^a} = e^{-i\omega_a \operatorname{tr} T^a} . \tag{6.21}$$

For the commutator $[\tau^a, \tau^b]$, we have

$$\mathrm{tr}\,[\tau^a, \tau^b] = \mathrm{tr}(\tau^a\tau^b - \tau^b\tau^a) \underset{\substack{\uparrow \\ \text{cyclicity}}}{=} 0 \qquad \text{trace free}$$

$$[\tau^a, \tau^b]^\dagger = [\tau^{b\dagger}, \tau^{a\dagger}] = [\tau^b, \tau^a] = -[\tau^a, \tau^b]$$

anti-hermitean

Hence, we can write $[\tau^a, \tau^b] = i h$ with h hermitean.
Since h can be spanned by τ^a again, we have

$$[\tau^a, \tau^b] = i f^{abc} \tau^c , \qquad (6.22)$$

where the f^{abc}'s are the structure constants of the Lie algebra $su(N_c)$ defined by (6.22).
Conventionally, the τ^a's are normalized to

$$\mathrm{tr}\, \tau^a \tau^b = \frac{1}{2} \delta^{ab} \qquad (6.23)$$

A well-known example is given by $N_c = 2$, where $\tau^a = \frac{1}{2}\sigma^a$ (Pauli matrices) such that

$$[\tau^a, \tau^b] = \frac{1}{4}[\sigma^a, \sigma^b] = \frac{1}{4} \cdot 2 i \varepsilon^{abc} \sigma^c$$
$$= i \varepsilon^{abc} \tau^c , \qquad (6.24)$$

i.e. the structure constants of $su(2)$ are $f^{abc}_{su(2)} = \varepsilon^{abc}$.

For all higher N_c, the generators can be constructed analogously to the Pauli matrices, e.g.

$N_c = 3$: $\quad N_c^2 - 1 = 8 \quad , \tau^a = \frac{1}{2}\lambda^a$

$$\lambda^1 = \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \lambda^2 = \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \lambda^3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \lambda^4 = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix},$$
$$\lambda^5 = \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, \lambda^6 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \lambda^7 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, \lambda^8 = \frac{1}{\sqrt{3}}\begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix} \quad (6.25)$$

These are the Gell-Mann matrices. The determination of the structure constants is straightforward:

$f^{abc}_{su(3)}$: (6.26)

123	147	156	246	257	345	367	458	678
1	$1/2$	$-\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{\sqrt{3}}{2}$	$\frac{\sqrt{3}}{2}$

and correspondingly for the permutations of the indices.

The resulting representation of $su(N_c)$ in terms of the τ^a are irreducible by construction. It is called the fundamental representation. Of course, higher representations of the same algebra (6.22), $[T^a, T^b] = i f^{abc} T^c$, in terms of higher dimensional matrices T^a also exist. An important one follows directly from the Jacobi identity for the commutator:

$$[[T^a, T^b], T^c] + [[T^b, T^c], T^a] + [[T^c, T^a], T^b] = 0$$

$$\Rightarrow f^{abd} f^{dce} + f^{bcd} f^{dae} + f^{cad} f^{dbe} = 0$$

$$\Rightarrow (-i f^{bad})(-i f^{edc}) - \underbrace{(-i f^{bcd})(-i f^{eda})}_{=(-i f^{ead})(-i f^{bdc})} = +i f^{bed} (-i f^{dac})$$

$$\Rightarrow [(-i f^b), (-i f^e)]^{ac} = i f^{bed} (-i f^d)^{ac} . \qquad (6.27)$$

Hence, $(T^a)^{bc} = -f^{abc}$ is also a representation of the $su(N_c)$ Lie algebra, consequentially generating a corresponding representation of $SU(N_c)$ in terms of $(N_c^2-1) \times (N_c^2-1)$ matrices. This is the <u>adjoint</u> repre<u>sentation</u>.

Now, let us start with a free Dirac theory for a massive quark field occuring in N_c colors:

$$\mathcal{L}_D = \bar{\psi}^i i\partial\!\!\!/ \psi^i - m \bar{\psi}^i \psi^i \quad , \quad i = 1, \dots, N_c \qquad (6.28)$$

As noted before, this theory is invariant under unitary global rotations in color space:

$$\psi^i \to U^{ij} \psi^j \quad , \quad \bar{\psi}^i \to \bar{\psi}^j (U^\dagger)^{ji} \qquad (6.29)$$

Using the representation (6.20), it is straightforward to show that the corresponding Noether current is given by

$$j^{\mu a} = \bar{\psi}^i \gamma^\mu T^a_{ij} \psi^j \quad , \quad \partial_\mu j^{\mu a} = 0 \quad (6.30)$$

Identifying $J^{\mu a} = -g\, j^{\mu a}$ with a coupling constant $g > 0$ as the vector-color current that we wish to couple to a photon-like color gauge field, we recognize that this color gauge field also has to carry an adjoint index:

in QED: $e < 0$

$$\mathcal{L}_J = -JA = -J^{\mu a} A^a_\mu \quad , \quad a = 1, \dots, N_c^2 - 1 \quad (6.31)$$

Adding the current term to the free Dirac theory, we obtain the Lagrangian

$$\mathcal{L} = \bar{\psi}^i \, i \not{D}_{ij} \psi^j - m \bar{\psi}^i \psi^i \qquad (6.32)$$

where the covariant derivative now takes the form

$$\not{D}_{ij} = \gamma^\mu D_{\mu\, ij} = \gamma^\mu \left(\partial_\mu \delta_{ij} - i g \tau^a_{ij} A^a_\mu \right) \qquad (6.33)$$

Incidentally, note that – in order to preserve to invariance of (6.32) under <u>global</u> color rotations – A^a_μ is not allowed to remain unmodified under a global rotation. Writing

$$A_{\mu\, ij} := \tau^a_{ij} A^a_\mu \qquad (6.34)$$

or A_μ in short, the color gauge field has to transform as

$$A_\mu \rightarrow U A_\mu U^\dagger \quad \text{under } \underline{\text{global}} \qquad (6.35)$$

color rotations. Note that this is still in line with QED, as for a $U(1)$ symmetry the generator is a number, say $\tau|_{U(1)} = 1$, such that $U A_\mu U^\dagger = A_\mu$ for QED.

However, inspired from QED we now wish to promote the invariance to a <u>local</u> invariance.

This is possible if the covariant derivative of the spinor transforms as

$$\not{D}\psi \rightarrow U(x)\not{D}\psi \qquad (6.36)$$

as analogously in QED, cf. (6.11) such that

$$\bar{\psi} i\not{D}\psi \rightarrow i\bar{\psi}\underbrace{U^{+}U}_{=\mathbb{1}}\not{D}\psi = \bar{\psi} i\not{D}\psi . \qquad (6.37)$$

This condition for the covariant derivative is met if we generalize (6.35) to the <u>local</u> transformation rule

$$A_\mu \rightarrow A_\mu' = U A_\mu U^{+} - \frac{i}{g}(\partial_\mu U) U^{+} \qquad (6.38)$$

Check:

$$\underline{D_\mu \psi} \rightarrow (\partial_\mu - ig A_\mu')\psi' = (\partial_\mu - ig\, U A_\mu U^{+} - (\partial_\mu U)U^{+}) U\psi$$

$$= U(\partial_\mu - ig A_\mu)\psi + \cancel{(\partial_\mu U)\psi} - \cancel{(\partial_\mu U)}\underbrace{U^{+}U}_{=1}\psi$$

$$= \underline{U D_\mu \psi} \qquad (6.39)$$

Having introduced a field A_μ that couples "photon-like" to the color charge of the quarks, we finally need to specify its dynamics by constructing a kinetic term for A_μ on the level of the action.

For this, we first note that the field strength in electrodynamics follows from the commutator of covariant derivatives,

$$U(1): \qquad [D_\mu, D_\nu] = ie F_{\mu\nu} \tag{6.40}$$

Taking the different sign conventions for the coupling into account, we similarly define the field strength for $SU(N_c)$ gauge theory from the covariant derivatives:

$$F_{\mu\nu} := \frac{1}{ig} [D_\mu, D_\nu] \qquad , F_{\mu\nu} \equiv F^a_{\mu\nu} t^a . \tag{6.41}$$

As discussed in the exercises, this leads to

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g f^{abc} A^b_\mu A^c_\nu . \tag{6.42}$$

As the covariant derivative transforms homogeneously,

$$D_\mu \to U D_\mu U^\dagger \qquad (\text{cf. } (6.36)) , \tag{6.43}$$

also the field strength transforms homogeneously,

$$F_{\mu\nu} \to U F_{\mu\nu} U^\dagger = F'_{\mu\nu} \qquad (6.44)$$

and is thus not invariant componentwise in contrast to electrodynamics.

Still, we can straightforwardly construct a gauge-invariant action

$$\underline{\underline{\mathcal{L}_{YM}}} = -\frac{1}{4} F^a_{\mu\nu} F^{a\mu\nu} \overset{(6.23)}{\equiv} -\frac{1}{2} F^a_{\mu\nu} F^{b\mu\nu} \operatorname{tr} \tau^a \tau^b$$

$$= -\frac{1}{2} \underline{\underline{\operatorname{tr} F_{\mu\nu} F^{\mu\nu}}} \qquad (6.45)$$

$$= -\frac{1}{2} \operatorname{tr} U^\dagger U F_{\mu\nu} U^\dagger U F^{\mu\nu}$$

$$\overset{\text{cyclicity}}{=} -\frac{1}{2} \operatorname{tr} (U F_{\mu\nu} U^\dagger)(U F^{\mu\nu} U^\dagger)$$

$$= -\frac{1}{2} \operatorname{tr} F'_{\mu\nu} F'^{\mu\nu} .$$

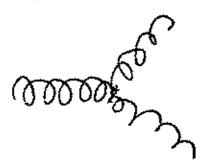
This is the celebrated Lagrangian of Yang-Mills theory, an $SU(N_c)$ bosonic theory of a vector field (spin-1) with a local symmetry.

It is important to realize that this action not only defines the kinetic terms for A^a_μ:

$$\mathcal{L}^{kin}_{YM} \simeq -\frac{1}{4} (\partial_\mu A^a_\nu - \partial_\nu A^a_\mu)(\partial^\mu A^{\nu a} - \partial^\nu A^{\mu a}) \qquad (6.46)$$

but also contains self-interaction terms which are enforced by gauge invariance (schematically)

$$\mathcal{L}_{YM}^{int} \sim + \dots \; g(\partial_\mu A_\nu) A^\mu A^\nu \quad + \dots \; g^2 (A_\mu A_\nu)^2 \qquad (6.47)$$



Therefore, already the pure Yang-Mills part is a highly non-trivial interacting theory unlike the pure Maxwell part. The gauge field excitations are also called gluons, hence (6.45) describes "gluodynamics". Read together with the Dirac part of the quarks (6.32), we arrive at the classical action defining Quantum Chromodynamics (QCD)

$$S_{QCD} = \int d^4x \left[-\frac{1}{4} F^a_{\mu\nu} F^{a\mu\nu} + \bar{\psi} i \not{D} \psi - m \bar{\psi} \psi \right]. \qquad (6.48)$$

Upon the inclusion of different quark flavors, each flavor may have a different mass parameter m.

Let us finally take a quick look at the classical field equations. As discussed in the exercises, the Euler-Lagrange equations for the gauge/gluon field lead to

$$D_\mu^{ab} F^{b\mu\nu} \equiv \left(\partial_\mu \delta^{ab} + g f^{acb} A_\mu^c \right) F^{b\mu\nu} = j^{a\nu} \qquad (7.49)$$

where $j^{a\nu} = \bar{\Psi} g \gamma^\nu T^a \Psi$. Here, we encounter the covariant derivative in the adjoint representation:

$$D_\mu^{ab} = \partial_\mu \delta^{ab} + g f^{acb} A_\mu^c = \left(\partial_\mu - i g T^c A_\mu^c \right)^{ab} \qquad (7.50)$$

with $(T^c)^{ab} = -i f^{cab}$.

Let us, for example, consider a static quark-anti-quark pair as a simple model for a meson,

$$j^{a0} = \underset{\substack{\uparrow \\ \text{charge}}}{Q} \; \underset{\substack{\uparrow \\ \text{unit vector} \\ \text{in color space} \\ n^2 = 1}}{n^a} \left(\delta^{(3)}(\vec{x} - \underset{\substack{\uparrow \\ \text{quark} \\ \text{position}}}{\vec{x}_1}) - \delta^{(3)}(\vec{x} - \underset{\substack{\uparrow \\ \text{antiquark} \\ \text{position}}}{\vec{x}_2}) \right) \qquad (7.51)$$

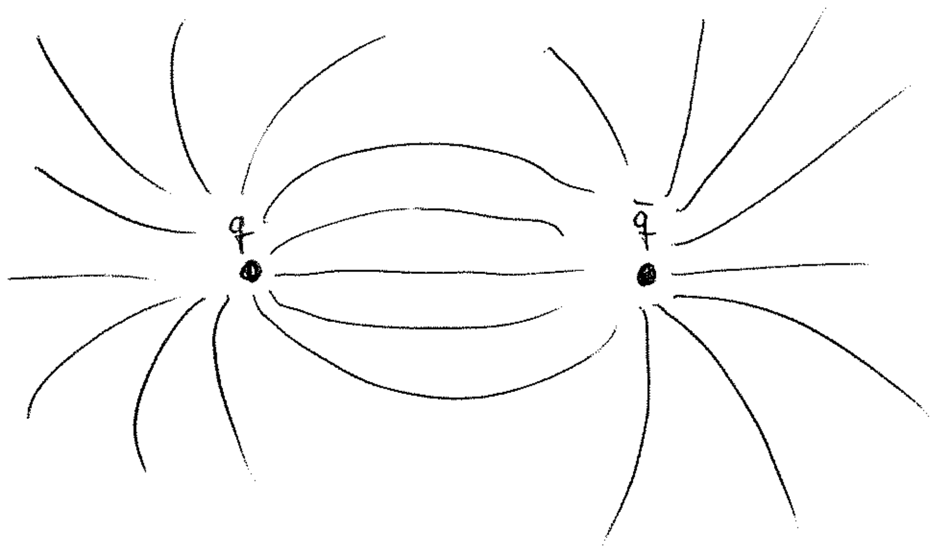
The equation of motion can fully be mapped onto classical electrodynamics, by noting that a pseudo-abelian ansatz

$$A^a_\mu = n^a a_\mu \quad , \quad F^a_{\mu\nu} = n^a f_{\mu\nu} \qquad (7.52)$$

with $f^{acb} n^b n^c = 0$ leads to

$$\partial_\mu f^{\mu\nu} = J^{a\nu} n^a . \qquad (7.53)$$

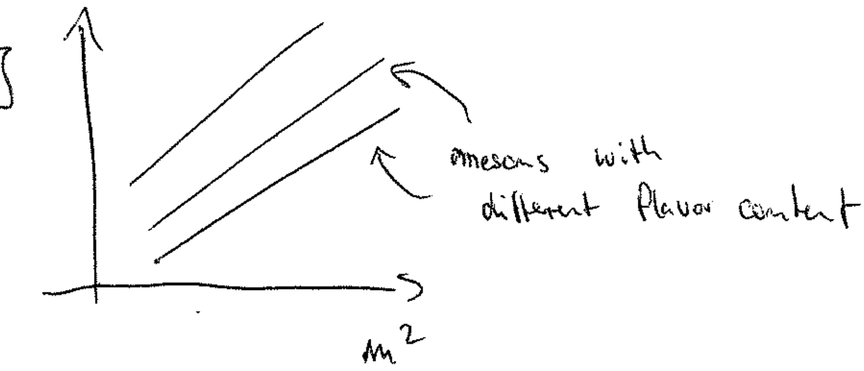
Hence, the solution is fully equivalent to that of a classical dipole field for the n^a component of the chromoelectric field:



Correspondingly, the static potential corresponds to the Coulomb potential

$$V(r) \sim \frac{1}{r} \qquad (7.54)$$

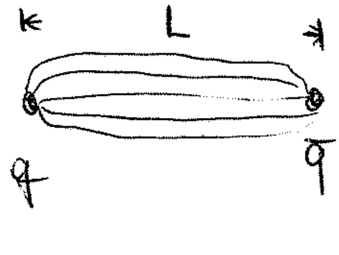
However, this is in contradiction with the experimental observation. For instance, if higher mesonic excitation with higher angular momentum J are studied, one observes that there total mass^2 is proportional to J: $\quad J \sim m^2 \qquad (7.55)$



These lines of proportionality are called "Regge trajectories".

In contrast to the classical analysis given above, this observation can be described by a string model for the field distribution of a meson:

- the gluon field of a meson is stringlike with a constant energy per length σ (string tension)



- For higher excitations, the quarks on both ends rotate at almost the speed of light

Then the energy/mass of the system is, $R = \frac{L}{2}$

$$m \simeq E = 2 \int_0^R \frac{\sigma}{\sqrt{1 - v^2_{(r)}}}\, dr = 2 \int_0^R \frac{\sigma \, dr}{\sqrt{1 - \left(\frac{r}{R}\right)^2}}$$

$$= \pi \sigma R \, , \tag{7.56}$$

whereas the angular momentum is

$$J = 2 \int_0^R \frac{\sigma \; r \, v(r) \, dr}{\sqrt{1 - v^2_{(r)}}} = \frac{2}{R} \sigma \int_0^R \frac{r^2 \, dr}{\sqrt{1 - \frac{r^2}{R^2}}}$$

$$= \frac{1}{2} \pi \sigma R^2 \, , \tag{7.57}$$

from which we read off that

$$J = \frac{1}{2\pi\sigma} \, m^2 \, . \tag{7.58}$$

This is in agreement with the experimental observation. The slope of the Regge trajectories gives

$$\alpha' = \frac{1}{2\pi\sigma} \simeq 0.9\ (\mathrm{GeV})^{-2}$$

$$\text{or} \quad \sigma \simeq (430\ \mathrm{MeV})^2 \qquad (7.59)$$

A stringlike color electric field distribution can be associated with a linear potential,

$$V(r) \sim r \qquad (7.60)$$

This line-like field distribution between two quarks and the corresponding impossibility to isolate a single quark is called „confinement".

The comparison with our conclusion from the classical equation of motion shows that classical Chromodynamics is insufficient to describe this basic experimentally verified property of the strong interactions.

Therefore: Quantum effects modify the dynamics of QCD qualitatively (not only quantitatively).