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On the connection between the scales of weak and strong interactions

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Abstract

We investigate within the standard model the possibility that nonperturbative QCD effects determine the Fermi scale and electroweak symmetry breaking is a consequence of chiral symmetry breaking. In this scenario the ratio between the Fermi scale and the quark condensate $\langle\bar{\psi}\psi\rangle_0^{1/3}$ comes out inversely proportional to the Yukawa coupling of the strange quark, consistent with observation. The Higgs particle mass is predicted in the range of 200 KeV.

The standard model of electroweak and strong interactions [1)2)] has two different mass scales: The Fermi scale φ_0 = 174 GeV determines the strength of weak interactions and the masses of quarks and leptons. It is given by the vacuum expectation value (vev) of the scalar doublet which spontaneously breaks weak SU(2) x U(1) symmetry. The QCD scale Λ_{QCD} (a few times hundred MeV) sets the scale for nuclear masses and interactions. It characterizes the scale where the gauge coupling g_s of SU(3) becomes strong. In addition, there are important features of the standard model which depend on a complicated interplay between φ_o and Λ_{QCD}. For example the pion mass m_π is proportional $(\varphi_o\Lambda_{QCD})^{1/2}$. It is obvious that even a moderate change in the ratio

$$\gamma = \frac{\Lambda_{QCD}}{\varphi_0} \approx 10^{-3} \qquad (1)$$

would lead to a very different picture in almost all branches of physics. (For example the electron to proton mass ratio is proportional γ^{-1}.)

It is certainly one of the most important challenges for a fundamental theory to explain why γ is of the order 10^{-3}. In addition, modern unification is often associated with a huge unification scale of the order of the Planck mass M_p. This not only leads to the puzzle how to understand the small ratio φ_o/M_p (gauge hierarchy problem [3)]) but also to the question why Λ_{QCD} and φ_o are so "near" to each other when looked upon from a characteristic scale M_p. From the viewpoint of the short distance physics around M_p the difference between Λ_{QCD} and φ_o appears like a "fine structure" in the effective long distance physics, similar in size to the structure in the fermion mass matrices [4)] reflected by different Yukawa couplings. We may call this the "connection problem" between the scales of weak and strong interactions. What has φ_o to do with Λ_{QCD}? The connection problem and the gauge hierarchy problem are of course not unrelated. Whenever the tiny ratio Λ_{QCD}/M_p is explained by the logarithmic evolution of the strong gauge coupling, a solution of the connection problem and thereby an understanding of γ would automatically solve the gauge hierarchy problem.

In perturbation theory the scales Λ_{QCD} and φ_0 are essentially unrelated free parameters of the standard model.[1] In presence of a scale Λ_{QCD} emerging from strong interactions, however, a perturbative treatment of electroweak symmetry breaking becomes questionable. Nonperturbative QCD effects lead to interactions between the σ field (quark-antiquark bound state) and the Higgs doublet of the type $\sigma^3\bar{\varphi}$, $\sigma^2\bar{\varphi}^2$ etc. The field $\bar{\varphi}$ corresponds to the mean value of the weak doublet in a volume V_{QCD} with characteristic length scale Λ_{QCD}^{-1}. These interactions are therefore local only for momenta below Λ_{QCD} but they become nonlocal when considered at length scales smaller than Λ_{QCD}^{-1}. It requires some thought to compare these nonlocal interactions with the local interactions described by the (classical) potential for the weak doublet.

To illustrate the problem consider first the pure Higgs model (four component φ^4 theory) in the spontaneously broken phase, such that the minimum of the perturbative potential V_p corresponds to some large scale M. (One may take M in the vicinity of the Planck scale.) As is well known [5] the effective scalar potential $\Gamma(\bar{\varphi})$ is convex. (It is the convex hull of the perturbative potential $V_p(\bar{\varphi})$.) In the broken phase perturbation theory can be trusted only for $|\bar{\varphi}| \geq M$ [2], but it breaks down for $|\bar{\varphi}| < M$ due to a failure of the saddle point approximation (compare fig. 1). This reflects the fact that only surface energy is needed to change in a large domain the phase of φ. Even though in every small volume with scale M^{-1} the magnitude of the

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1) There is a possible exception for the case of seven or eight generations where the strong gauge coupling increases substantially only for momenta below the Fermi scale.

2) There is convincing evidence [6] for the reliability of perturbation theory even for the case of strong quartic bare couplings.

mean value of φ is near M, the probability of finding configurations with mean values $\bar{\varphi} = 0$ or $\bar{\varphi} = M$ becomes almost degenerate for large volumes.

Let us now include strong interactions and quarks with Yukawa couplings. For a first crude estimate we suppose that the "electroweak" part of the potential relevant at the scale Λ_{QCD} is well approximated by $\Gamma(\bar{\varphi})$. We simply add the nonperturbative QCD effects which will act as perturbations on $\Gamma(\bar{\varphi})$. A perturbation linear in $\bar{\varphi}$ pushes the minimum of $\Gamma(\bar{\varphi}) + \tilde{\alpha}\sigma^3\bar{\varphi}$ to a now uniquely selected minimum of V_p at $|\bar{\varphi}| = M$. Here perturbation theory applies and the nonperturbative QCD effects are completely negligible. This situation can change drastically for nonlinear perturbations. It becomes possible that the nonlinear perturbation $V_{\varphi\sigma}(\sigma, \bar{\varphi})$ develops a minimum within the flat region of $\Gamma(\bar{\varphi})$. The minimum of $\Gamma + V_{\varphi\sigma}$ and therefore the expectation value of $\bar{\varphi}$ is then determined by the minimum of $V_{\varphi\sigma}$ and not by V_p! Intuitively the QCD effects can favour energetically a certain mean value of $\bar{\varphi}$ within a volume V_{QCD}. Since V_{QCD} is large compared to a volume with length scale M^{-1} it costs very little "electroweak" energy to arrange the domains within V_{QCD} so that this mean value obtains. The nonperturbative QCD effects could dominate the effective potential!

At long distances strong interactions can be described by an effective (linear) σ model. Quark-antiquark pairs $q_L\bar{q}_R$ form scalar mesons which transform as doublets under SU(2) with hypercharge one - just the same as the weak doublet φ. Chiral symmetry is spontaneously broken by a vev $\sigma_0 \approx \frac{1}{2} f_\pi$ = 67 MeV.[3] Due to its Yukawa couplings to quarks the mean value $\bar{\varphi}$ will interact with the $\bar{q}q$ condensate and therefore with σ. The interaction linear in $\bar{\varphi}$ has the form $(\tilde{\sigma} = i\tau_2\sigma^*)$

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3) In absence of φ the W and Z bosons would acquire a mass of the order f_π.

$$\mathcal{L}^{(1)}_{\varphi\sigma} = \alpha(\sigma^{+}\sigma)(\sigma^{+}\bar{\varphi}) + \alpha'(\sigma^{+}\tilde{\sigma})(\tilde{\sigma}^{+}\bar{\varphi}) + h.c. \qquad (2)$$

This has three immediate consequences: First, there will be a vacuum alignment between $\bar{\varphi}$ and σ. If we choose a convention where the lower component of σ has a real vev σ_o, we find that a vev of the lower component of $\bar{\varphi}$ is energetically favoured compared to the upper component. This correlation guarantees that the electromagnetic U(1) symmetry remains unbroken. Similarly, the phase of $\langle\bar{\varphi}\rangle = \varphi_o$ (compared to σ_o) is dictated by the phases of α and α'. The correlation between the phases of σ_o and φ_o may have implications for the CP problem.

Second, effective terms involving $(\sigma^{+}\bar{\varphi})$ break the global $SU(2N_G)_L \times SU(2N_G)_R$ flavour symmetry which would exist in the absence of electroweak interactions. The vev of $\bar{\varphi}$ induces masses for the quarks and for $\varphi_o \neq 0$ the pions (and similarly other mesons) acquire a mass.

Finally, the interaction (2) puts a lower bound on the ratio φ_o/σ_o if weak interactions are in the spontaneously broken phase. (By this we mean in our context that the potential for $\bar{\varphi}$ - neglecting its interactions with σ - should not contain a positive quadratic term.) A potential of the form $V = -\tilde{\alpha}\sigma_o^3\bar{\varphi} + \frac{1}{2}\lambda_\varphi\bar{\varphi}^4$ leads to

$$\varphi_o = \left(\frac{\tilde{\alpha}}{2\lambda_\varphi}\right)^{1/3}\sigma_o \qquad (3)$$

and a negative quadratic term $-\mu_\varphi^2\bar{\varphi}^2$ only increases φ_o/σ_o. Additional interactions with σ influence the quantitative value for φ_o/σ_o but do not change the order of magnitude of the bound. Without the nonperturbative QCD effects bounds for the Higgs mass can be obtained for a given value of φ_o [8)6)], but a lower bound on the vev φ_o itself does not exist unless the term (2) is included. All these effects of the linear term are independent of the detailed properties of the electroweak potential.

For the observed value of φ_o the interaction between σ and $\bar{\varphi}$ is dominated by the Yukawa coupling of the strange quark. We estimate $\tilde{\alpha} \approx 30\, h_s$ using the identification

$$V_{\varphi\sigma} = -\tilde{\alpha}\sigma_o^3\bar{\varphi} \quad + \text{terms nonlinear in } \bar{\varphi} + \text{const} \qquad (4)$$

$$\approx h_u\bar{\varphi}\langle\bar{u}u\rangle + h_d\bar{\varphi}\langle\bar{d}d\rangle + h_s\bar{\varphi}\langle\bar{s}s\rangle + h_c\bar{\varphi}\langle\bar{c}c\rangle + \cdots$$

The contribution of up and down quarks can be neglected and the condensates of heavier quarks are in leading order inversely proportional to the quark mass [9)]. Alternatively, we can extract $\tilde{\alpha}$ (and also terms nonlinear in $\bar{\varphi}$) from the phenomenological analysis of chiral perturbation theory [10)]. (The term $\sim\tilde{\alpha}\sigma_o^3\bar{\varphi}$ leads to the quark mass term in the nonlinear σ model.)

There is no reason why the interactions between σ and $\bar{\varphi}$ should be linear in $\bar{\varphi}$. Also the $\bar{\psi}\psi$ condensates depend themselves on $\bar{\varphi}$ and this introduces an effective nonlinearity in $V_{\varphi\sigma}$. Let us concentrate on the term $h_s\bar{\varphi}\langle\bar{s}s\rangle$. We insert the vacuum value $\sigma_o(\bar{\varphi})$ and express the resulting $\bar{\varphi}$ dependence of $V_{\varphi\sigma}$ through the dependence of the condensate $\langle\bar{s}s\rangle$ on the strange quark mass m_s

$$\frac{dV_{\varphi\sigma}(\bar{\varphi})}{d\bar{\varphi}} = h_s\frac{d}{dm_s}\Big(m_s\langle\bar{s}s\rangle(m_s)\Big) \qquad (5)$$

For low values of m_s the strange quark is effectively massless and $\langle\bar{s}s\rangle$ becomes independent of m_s. Extrapolating the QCD sum rule estimate [9)] for large m_s

$$\langle\bar{s}s\rangle = -\frac{1}{12m_s}\left\langle\frac{\alpha_s}{\pi}G^a_{\mu\nu}G_a^{\mu\nu}\right\rangle \qquad (6)$$

to a strange quark mass of the order 150-200 MeV gives a value for $\langle\bar{s}s\rangle$ which is a factor 2-3 smaller in size than $\langle\bar{\psi}\psi\rangle_o \simeq (230\text{ MeV})^3$. This suggests that (6) is approached from below and therefore $m_s\langle\bar{s}s\rangle(m_s)$ should have a minimum at $\bar{m}_s$. The scale for $\bar{m}_s$ is obviously given by $\langle\bar{\psi}\psi\rangle_o^{1/3}$. This is in the vicinity of the physical strange quark mass m_s^{phys}. The same conclusion is suggested by an expansion in m_s within chiral perturbation theory [10]. Taking $\varepsilon_s = 0.3$ in ref. 10 leads to a minimum value $\bar{m}_s \approx m_s^{phys}$. The minimum of $V_{\varphi\sigma}(\bar{\varphi})$ corresponds to a value $\bar{m}_s$ where the strange quark changes its role from a light to a heavy quark, $\varphi_o = \bar{m}_s/h_s$. It is puzzling that our naive estimate of φ_o fits well the observed value $\varphi_o = m_s^{phys}/h_s = 174$ GeV!

For a more accurate treatment we have to account for the fact that Λ_{QCD} (and therefore the gluon condensate and $\langle\bar{\psi}\psi\rangle_o$) depends on $\bar{\varphi}$ via the heavy fermion mass thresholds in the evolution equation for the strong gauge coupling. Also the other quarks have to be included in $V_{\varphi\sigma}$. It is still plausible that $\bar{\varphi} = m_s/h_s$ is near a local minimum of $V_{\varphi\sigma}$. The question if it is a global minimum seems more complicated. The quadratic term at the minimum [4], $M_H^2 = \frac{1}{2} d^2 V_{\varphi\sigma}/d\bar{\varphi}^2(\varphi_o)$, depends on the smoothness of the transition between light quark and heavy quark regime. For a rough estimate we approximate

$$V_{\varphi\sigma} \approx -h_s \langle\bar{\psi}\psi\rangle_o \bar{\varphi} + \beta \langle\bar{\psi}\psi\rangle_o^{2/3} \bar{\varphi}^2$$

$$\beta = h_s \langle\bar{\psi}\psi\rangle_o^{1/3} / 2\varphi_o = O(h_s^2) \tag{7}$$

and obtain

$$M_H^2 = \frac{h_s \langle\bar{\psi}\psi\rangle_o}{2\varphi_o} \approx (0.2\ \text{MeV})^2 \tag{8}$$

This is an unusually small value for the Higgs particle mass! It is easy to understand the large value for φ_o/σ_o intuitively: Although the linear term driving $\bar{\varphi}$ away from the origin at $\bar{\varphi} = 0$ is small $\sim h_s$, a minimum can only occur once the restoring force is of equal strength than the driving term. Restoring forces are suppressed by even higher powers of h_s ($\beta \sim h_s^2$). Therefore φ_o must come out proportional to the inverse of the small coupling h_s.

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4) Strictly speaking there is a mass matrix involving σ and $\bar{\varphi}$. Corrections to the eigenvalues are small ($O(h_s)$) in the present case.

For our estimate of φ_o we have made a very crude assumption, namely that the relevant scalar potential is $U(\bar{\varphi}) = \Gamma(\bar{\varphi}) + V_{\varphi\sigma}(\bar{\varphi})$. The effective potential Γ is relevant for an infinite volume whereas we are interested in the volume V_{QCD}. Also the short distance fluctuations of QCD have to be incorporated. (For a most naive guess they only lead to a modification of the perturbative potential V_p and in consequence to a shift in the value of M.) This suggests the use a "finite volume potential" $U(\bar{\varphi})$ defined symbolically (in Euclidean space) by

$$U_k(\bar{\varphi}) = -\frac{1}{\Omega} \ln \int_\Omega \mathcal{D}\varphi\, \delta\left(\frac{1}{V_k}\int_{V_k} dx\, \varphi(x) - \bar{\varphi}\right) \exp - S_\Omega[\varphi] \tag{9}$$

The exponential $\exp(-\Omega U_k(\bar{\varphi}))$ is a measure for the probability that the mean value of φ within a finite volume $V_k \sim k^{-d}$ is given by $\bar{\varphi}$. (The volume Ω is used as a purely technical device and the limit $\Omega \to \infty$ should be taken at the end.) The finite volume potential interpolates between $\Gamma(\bar{\varphi})$ and the "tree potential" $V_o(\bar{\varphi})$ defined at M

$$\lim_{k\to 0} U_k(\bar{\varphi}) = \Gamma(\bar{\varphi}) \quad , \quad \lim_{k\to M} U_k(\bar{\varphi}) = V_o(\bar{\varphi}) \tag{10}$$

It is not necessarily convex for nonvanishing k. (This potential is similar - up to effects from the boundary of V_k - to the "constraint effective potential" [11]). We may interprete the definition (9) as a specific way of integrating out fluctuations in φ with typical length scales smaller than k^{-1}. Choosing a coordinate y in the center of V_k we can consider $\bar{\varphi}(y)$ as a "block spin" variable. Writing the Euclidean partition function

$$Z = \int_\Omega D\varphi \, \exp - S[\varphi]$$

$$= \int_\Omega D\bar{\varphi}(y) \int_\Omega D\varphi(x)\, \delta\left(\frac{1}{V}\int_V \varphi dx - \bar{\varphi}(y)\right) \exp - S[\varphi] \qquad (11)$$

$$= \int_\Omega D\bar{\varphi}(y) \exp\left(-\int_\Omega dy\, \mathcal{L}(\bar{\varphi}, \partial_\mu \bar{\varphi}, \ldots)\right)$$

we identify $U(\bar{\varphi})$ with $\mathcal{L}(\bar{\varphi})$ for constant $\bar{\varphi}$. (One should use a renormalization of the variable $\bar{\varphi}(y)$ to bring the kinetic term in $\mathcal{L}$ into standard form. The δ distribution could also be replaced by an appropriate Gaussian.) The short distance (perturbative) QCD contributions can be incorporated into $U_k(\bar{\varphi})$ by an appropriate extension of the block spin approach. The nonperturbative effects from chiral symmetry breaking can now be included by adding $V_{\varphi\sigma}(\bar{\varphi})$ to $U_k(\bar{\varphi})$ at a scale $k \approx \Lambda_{QCD}$. (Note that $U_{QCD}(\bar{\varphi})$ is also a function of the mean value of the weak doublet in a volume V_{QCD}.) The combined potential $U + V_{\varphi\sigma}$ can be extrapolated to scales below Λ_{QCD} and approaches the full effective potential for $k \to 0$. If the minimum of $U_{QCD} + V_{\varphi\sigma}$ is determined essentially by $V_{\varphi\sigma}$ the flat region of the full effective potential extends only between the minima of $V_{\varphi\sigma}$ istead of V_p. The importance of $V_{\varphi\sigma}$ depends on how fast $U_k(\bar{\varphi})$ converges to $\Gamma(\bar{\varphi})$ in the spontaneously broken phase. Since the difference between the two quantities involves only surface effects it is plausible that $\Gamma(\bar{\varphi})$ becomes a good approximation to $U_k(\bar{\varphi})$ at $k \approx \Lambda_{QCD}$ for a wide range of "tree" parameters $\mu^2_\varphi(M)$ and $\lambda_\varphi(M)$, including mass terms $\mu^2_\varphi(M)$ of the order of M^2 which would lead to a minimum of the perturbative potential at a large scale M.

Leptons and their Yukawa couplings are easily added. Inclusion of the electroweak gauge interactions needs an even more complex discussion since the mean value $V^{-1}\int_V dx\, \varphi(x)$ is not a gauge invariant quantity anymore. There are no massless Goldstone bosons in this case. Nevertheless, the flatness of quantities

like $U(\bar{\varphi})$ is not directly related to the existence of propagating massless modes but only to the fact that the coherence in the phase of the mean value of φ becomes weak over large distances. In a gauge fixed version the problem looks at first sight not very different from the global SU(2) x U(1) model discussed before.

Two possible objections against our scenario should be shortly addressed. One concerns the stability of our approach with respect to radiative corrections. One may wonder if fermion or gauge boson loops with momenta near Λ_{QCD} or φ_0 do not induce terms $\mu^2 \bar{\varphi}^2$ or $\lambda \bar{\varphi}^4$ compared to which the nonperturbative QCD effects are completely negligible. Such loops induce effective nonlocal scalar interactions[5] and they influence the way how $U_k(\bar{\varphi})$ approaches $\Gamma(\bar{\varphi})$. The behaviour of $U_k(\bar{\varphi})$, however, cannot be determined by standard perturbation theory. (The standard perturbative renormalization group equations for μ^2 and λ are inadequate for $|\bar{\varphi}| < M$. In contrast, the renormalization group analysis remains valid for Yukawa couplings and gauge couplings.) An appropriate (perturbative?) block spin formalism is needed. This must account for the fact that the behaviour of $U_k(\bar{\varphi})$ is determined by the physics of domain adjustment in the presence of light fermion and gauge boson loops. Not much about the importance of these effects can be said before.

The other is about the relevance of the QCD σ-model for weak processes involving the exchange of W or Z bosons. Naively one could think that at momenta around the Fermi scale the nonperturbative QCD effects should not be relevant. The weak bosons propagate through the vacuum for which $\bar{\varphi} = \varphi_0$. The vev φ_0 sets the scale for all weak interaction processes. Once we know φ_0 we can forget its origin, use perturbative weak interaction field theory and neglect all nonperturbative QCD effects. The value of φ_0 itself, however, is a property of the surrounding

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5) In principle the combination of such effects could also produce a nontrivial minimum of U_k and determine φ_0.

vacuum and related to zero momentum[6] rather than to the momenta of a particular weak scattering process. Its value is therefore sensitive to the long distance behaviour of the theory where nonperturbative QCD effects become important.

Many questions are left open. Nevertheless it seems not excluded that electroweak symmetry breaking is indeed determined by the chiral symmetry breaking in QCD. The connection between nonperturbative QCD effects and (perturbative) electroweak physics certainly needs and merits a more profound investigation. This shows that there are still important holes in our understanding of spontaneous symmetry breaking within the standard model. Although we are aware of the somewhat speculative character of this hypothesis it seems worthwhile to mention some of the consequences of this version of electroweak symmetry breaking: The gauge hierarchy problem would be solved and φ_0 becomes calculable in terms of Λ_{QCD} and Yukawa couplings. There is a possible reduction in the number of parameters since the parameters $\mu_\varphi^2(M^2)$ and $\lambda_\varphi(M^2)$ become irrelevant if they correspond to a region in parameter space for which $U_{QCD}(\bar\varphi)$ is effectively flat. This could also have important consequences for the issues of dilatation symmetry, the cosmological constant and a possible new intermediate range interaction [12] since Λ_{QCD} is now the only low energy mass scale. [7] The mass of the Higgs scalar is predicted to be very low ($\approx$ 200 keV). The scalar is therefore not expected to decay into e^+e^- and must have a rather long lifetime. A scalar of this type is not excluded experimentally so far. It may be possible to detect it by future precision experiments. There are experimental bounds [13] on the scalar couplings to nucleons. These are not easy to evaluate

6) One may ask if for a weak process the relevant quantity is not the mean value of the doublet in a volume of Fermi size (φ_0^{-1}). In any case, by simple dimensional arguments, this mean value cannot be very different from φ_0 .

7) If there is a cosmon force its range should be in the kilometer range ($\sim M_p/\Lambda^2_{QCD}$).

theoretically. In particular we should mention that the Higgs scalar has residual "strong" interactions since there is a mixing with the σ field of order $\tilde{\alpha} \sim h_s$. The physical Higgs scalar is a sort of collective excitation. At energies below Λ_{QCD} it can be treated as a fundamental scalar, but higher energy scales need a detailed investigation. The chiral and weak phase transitions would look quite different from the standard picture. This has possibly important consequences for the early universe. In view of these prospects the questions concerning the connection between the scales of strong and electroweak interactions should be an important task for a deeper field theoretical investigation of symmetry breaking in the standard model.

Figure caption

Effective potential $\Gamma(\bar\varphi)$ and perturbative potential $V_p(\bar\varphi)$ in the spontaneously broken phase.

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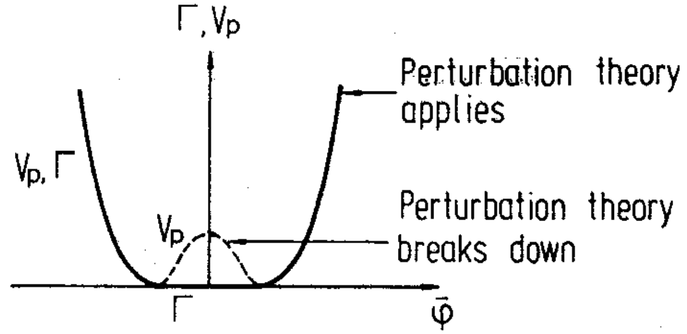

Γ, Vp
Perturbation theory applies
Vp, Γ
Perturbation theory breaks down
Vp
Γ
φ̄

Fig. 1