Lattice simulations of the strange quark mass and Fritzsch texture

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Abstract

A number of numerical simulations of lattice gauge theory have indicated a low mass of strange quark in 100 MeV range at the scale of $\mu = 2$ GeV. In the unquenched case, which is improved over the simulation in the quenched approximation by the inclusion of $u$ and $d$ sea quark effects, one sees a further downward trend. Here the fermion mass spectrum of the Fritzsch texture is recalculated. In a single step supersymmetric GUT with $M_X \sim 10^{16}$ GeV such values of the strange quark mass can be obtained for low values of $\tan \beta$. Experimental numbers $m_t^{pole} = 173 \pm 6$ GeV and $4.1 < m_b(m_b) < 4.4$ GeV are used in this study. Since the scenario is supersymmetric, gaugino loop diagrams contribute to the masses in addition to usual tree level Yukawa contributions. Upper bound of the mixing parameter $V_{cb}$ is taken at 0.045.
The strange quark mass has been traditionally calculated using the current algebra mass ratio\[1\]

\[
\frac{m_s}{m_u + m_d} = 12.6 \pm 0.5
\]  

(1)

Equation (1) is evaluated using values of \((m_u + m_d)\) which are reported in calculations of QCD finite energy sum-rules (FESR). At the two-loop level of perturbative QCD calculations which include non-perturbative corrections up to dimension six one has the result\[2\]

\[(m_u + m_d)(1 \text{ GeV}) = 15.5 \pm 2.0 \text{ MeV}\]  

(2)

For \(\alpha_s(m_Z) = 0.118\) results of (1) and (2) together lead to

\[m_s(1 \text{ GeV}) = 195 \pm 28 \text{ MeV} \quad \text{or} \quad m_s(2 \text{ GeV}) = 150 \pm 21 \text{ MeV}.\]  

(3)

The ratio

\[m_s(2 \text{ GeV})/m_s(1 \text{ GeV}) = 0.769 \quad \text{for} \quad \alpha_s(m_Z) = 0.118\]  

(4)

can be obtained by solving renormalization group equations\[3\]. A systematic uncertainty in this result remains in reconstruction of so called ‘spectral function’ from experimental data of resonances. When a different functional form of the resonance is adopted, and three loop order perturbative QCD theory is used one obtains\[4\]

\[(m_u + m_d)(1 \text{ GeV}) = 12.0 \pm 2.5 \text{ MeV}\]  

(5)

With (5) and (1) one gets

\[m_s(1\text{GeV}) = 151 \pm 32 \text{ MeV} \quad \text{or} \quad m_s(2\text{GeV}) = 116 \pm 24 \text{ MeV}.\]  

(6)

Again this translation from the scale of 1 GeV to the scale of 2 GeV is obtained for the case \(\alpha_s = 0.118\). It has been remarked in Ref.\[5\] that it is indeed difficult to account for vacuum fluctuations, or sea quark effects generated by quarks of small masses in perturbative QCD calculations. Thus, numerical simulations of strange quark mass on a lattice becomes rather attractive, especially if the simulation includes virtual light quark loop effects.
Up and down type quarks differ only in the $U(1)_{em}$ quantum numbers in an effective theory where the gauge symmetry is $SU(3)_c \times U(1)_{em}$. Lattice calculations in current literature have neglected effects of $U(1)_{em}$ which distinguishes up quarks from down quarks. Let us note that we are describing the lattice in terms of a theory at the scale of a few GeVs where light quark masses are to be described in terms of observables relevant to their own scales, which are meson masses and decay constants. Thus, lattice simulation determines $m_s, \frac{m_u + m_d}{2}$ and the lattice spacing $a$ using three hadronic observables. They can be chosen, for example, $M_\pi, M_{K^*}$, and $f_\pi$. Due to the structure of equations which are to be fitted, the scale $a$ can also be taken as a function of some other observable, for example, it may be chosen as $a(M_n)$ or $a(M_\Delta)$ etc. The result depend on the choice of the observable that fits the lattice spacing. The best choice would be the one which has minimum experimental uncertainty and the best result would be a clever weighted average of results from various choices. A test of the simulation is obviously to see whether results from various observables are statistically consistent with each other.

Next question is how do we describe quark masses when the theory is living on a discretized lattice. Various definitions or formalisms of quark masses on a lattice have been suggested. Ref.\cite{6} uses the definition in terms of hopping parameter $\kappa$ of the lattice

$$a \ m_{bare} = \log \left( 1 + \frac{1}{2\kappa} - \frac{1}{2\kappa_c} \right),$$

for a Wilson-like fermion. In the continuum limit we have $a \to 0$, and there one gets the hopping parameter $\kappa = \kappa_c = 1/8$. A smaller hopping parameter makes the lattice more sticky, and fermions remain on lattice points for a longer time. This make them look as if they were more massive.

There are various other formalisms of defining the mass of fermions on the lattice such as staggered fermions or domain wall fermions. Most calculations, however, use the Wilson action for various definitions of the fermion mass. Next, to compare the result with experiment, one has to calculate the \(\overline{MS}\) mass at a scale $\mu$ starting from the lattice estimate of the bare mass (7) using, for example, the mass renormalization constant $Z_m(\mu)$ relating the lattice regularization scheme to the continuum regularization scheme. The lattice regularization prescription is given in Ref\cite{7}. The $Z_m$ constant for various formalisms such as Wilson-like or Staggered are given in table\cite{4} of
Final results of the physical quark mass for various definitions of the fermion on a lattice differ $O(a)$ among each other and one expects to get the same result of the physical quark mass in the continuum limit when $a \to 0$.

Beyond the minimal lattice simulation of light quark masses using the heavy quark effective theory, the next step would be to incorporate sea quark effects. From the conservation of energy it can be understood that it is easiest to produce lightest of the quarks virtually. Indeed such simulations have been performed. They are termed $n_f = 2$ unquenched lattice simulations. The detailed processes of numerical simulations are described by respective authors. However, we have summarized the results of recent studies are in table 1 in the order: (A)[8], (B)[9], (C)[10], (D)[11], (E)[6], (F)[12], and (G)[13].

| reference | quenched | dynamical | (1/a)calibration | $m_s$(MeV) |
|-----------|----------|-----------|-----------------|----------|
| A         | yes      |           | $m_\rho$        | $143 \pm 6 & 115 \pm 2$ |
| B         | yes      |           | $m_\rho$, $m_{K^*}$ | $130 \pm 20$ |
| C         | yes      |           | $m_{K^*}$       | $122 \pm 20$ |
| D         | yes      |           | $m_{K^*}$       | $111 \pm 12$ |
| E         | yes      |           | $m_\rho$        | $110 \pm 31$ |
| F         | yes      |           | $m_\rho$        | $108 \pm 4$ |
| G         | yes      |           | $1P - 1S$ splitting | $95 \pm 16$ |
| A         | yes      |           | $m_\rho$        | $70 & 80$ |
| E         | yes      |           | $m_\rho$        | $68 \pm 19$ |
| G         | yes      |           | $1P - 1S$ splitting | $54 - 92$ |

Table 1: Ref. G uses 1P-1S splitting of the charmonium system to calibrate $(1/a)$. Reference A quotes two different results for two sets of parameters. All results are at the scale $\mu = 2$ GeV.

On the experimental front the bottom quark mass is in the range

$$4.1 < m_b(m_b) < 4.4 \text{ GeV}$$

(8)

according to the review of particle physics (PDG) tables[14]. Theoretically, one re-expresses the bottom quark mass in terms of parameters of the Minimal Supersymmetric Standard Model(MSSM). The tree level contribution
which is related straight to the Yukawa texture, and the one loop contribution due to the dominant gaugino loop can be accounted individually. Then one can write down the relation

\[ m_b = m_{\text{texture}}^{\text{texture}} + m_b^{\text{SUSY}} \]

\[ = h_b \frac{V_F}{\sqrt{2}} \cos \beta + m_b \frac{8}{3} g_3^2 \tan \beta \frac{m_\tilde{g} \mu}{16 \pi^2} \frac{m_{\tilde{g}}^2}{m_{\text{eff}}^2}. \]

(9)

Here \( m_\tilde{g} \) is the gluino mass \( \mu \) is the \( \mu \) parameter and \( m_{\text{eff}} \) is averaged supersymmetry breaking mass scale. This paper discusses a scenario where the first term of the RHS of (9) comes from diagonalizing a Fritzsch Yukawa texture. The second term can be estimated to be around \( \pm 2 \) GeV. In the supersymmetric case it will be satisfactory if the Fritzsch Yukawa contribution is in the range \( 2.1 < m_b^{\text{texture}} < 6.4 \) GeV

(10)

Next question is concerning \( \tan \beta \). Supersymmetry, together with the gauge quantum number structure of fermions demands that at least two Higgs doublets are necessary. Hence the ratio of the VEVs of two Higgs doublets is an unavoidable parameter given the value of the effective four-Fermi coupling \( V_F \). There are perturbative bounds on \( \tan \beta \) in the context of grand unified theories (They can be extended to supersymmetric theories without grand unification if MSSM is valid up to a certain high scale, say \( 10^{19} \) GeV). In practice the very low-valued regions of the parameter space for \( \tan \beta \) are forbidden from perturbative considerations. See [16] for example. Furthermore high values of \( \tan \beta \) have constraints from charge and color breaking[17, 18]. This is especially true if Yukawa couplings are at the fixed point region. The intermediate regions of \( \tan \beta \) are definitely allowed. To make a safe case let us choose the range for the purpose of this paper

\[ \tan \beta = 2 - 30. \]

(11)

Now let us focus on the texture. It has been noted that the quark mixing angle \( V_{us} \), which is a dimension-less quantity, can be thought of as a ratio of the mass scales of flavor symmetry breaking. These symmetries lead to the mass hierarchy between families. Phenomenologically of course, the ratio of the masses of the first and the second generation satisfies well the relation

\[ \tan \theta_c = \sqrt{\frac{m_d}{m_s}}. \]

(12)
If there are two Higgs doublets instead, (12) remains untouched as the ratio of the VEVs of the doublets cancel in the ratio on the RHS. Thus it cannot feel $\tan \beta$.

Suppose in a two generation case rotation angles of the up and the down sectors are $\theta_u$ and $\theta_d$. Then the combined quark mixing matrix $V = O_u O_d^\dagger$ will give $\theta_c = \theta_u \pm \theta_d$. This observation plays a role in the Fritzsch mass matrices. Fritzsch mass matrices can be thought of as a set of mass matrices which generalizes (12) to the following form\cite{19,20}

$$\theta_c = \theta_d \pm \theta_u = \tan^{-1} \sqrt{\frac{m_d}{m_s}} \pm \tan^{-1} \sqrt{\frac{m_u}{m_c}}.$$ (13)

Where $m_i$ are eigenvalues of Fritzsch mass matrices. In the three generation case Fritzsch textures for up and down sectors are given by

$$M_U = \begin{pmatrix} 0 & a e^{i\tau} & 0 \\ a e^{i\tau'} & 0 & b e^{i\sigma} \\ 0 & b e^{i\sigma'} & c e^{i\rho} \end{pmatrix} \quad M_D = \begin{pmatrix} 0 & A e^{i\rho'} & 0 \\ A e^{i\rho} & 0 & B e^{i\nu} \\ 0 & B e^{i\nu'} & C e^{i\omega} \end{pmatrix}. \quad (14)$$

The phases $r, R, r', R', h, H, h', H', q, Q$ can be absorbed in the redefinition of quark fields and individual mass matrices can be made real. However the weak charge changing current consists of (couples to) gauge eigenstates. Furthermore the structure of the weak charge changing current is $\psi_L^U \gamma^\mu \psi_L^D$. Consequently the weak mixing matrix must contain some combination of phases. It can be shown that residual phases of the weak mixing matrix are

$$\sigma = (r - R) - (h - H) - (h' - H') + (q - Q)$$

$$\tau = (r - R) - (h' - H').$$ (15)

when the weak mixing matrix is expressed as

$$O_U \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\sigma} & 0 \\ 0 & 0 & e^{i\tau} \end{pmatrix} O_D^{-1}. \quad (16)$$

In (16) $O_U$ and $O_D$ diagonalizes $\mathcal{M}_U$ and $\mathcal{M}_D$ which are precisely those in (14) but in the limit when all the phases vanish. Which are very simply

$$\mathcal{M}_U = \begin{pmatrix} 0 & a & 0 \\ a & 0 & b \\ 0 & b & c \end{pmatrix} \quad ; \quad \mathcal{M}_D = \begin{pmatrix} 0 & A & 0 \\ A & 0 & B \\ 0 & B & C \end{pmatrix}. \quad (17)$$
Let us consider the limit $m_u << m_c << m_t$ and $m_d << m_s << m_b$. We can approximately re-write (17) as following

$$M_U = \begin{pmatrix} 0 & a & 0 \\ a & -m_c & 0 \\ 0 & 0 & m_t \end{pmatrix}; \quad M_D = \begin{pmatrix} 0 & A & 0 \\ A & -m_d & 0 \\ 0 & 0 & m_t \end{pmatrix}. \tag{18}$$

Now we easily see that $M_U$ can be diagonalized by the rotation

$$O_U = \begin{pmatrix} \cos \theta_u^1 & \sin \theta_u^1 & 0 \\ -\sin \theta_u^1 & \cos \theta_u^1 & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \theta_u^2 & \sin \theta_u^2 \\ 0 & -\sin \theta_u^2 & \cos \theta_u^2 \end{pmatrix} \tag{19}$$

where $\tan \theta_u^1 = \sqrt{\frac{m_u}{m_c}}$ and $\tan \theta_u^2 = \sqrt{\frac{m_u}{m_t}}$. $M_D$ can be diagonalized similarly.

In the limit of strongly hierarchical eigenvalues one may use the approximation

$$\cos \theta_u^i \sim 1 \quad \cos \theta_d^i \sim 1. \tag{20}$$

Furthermore let us define the quantities

$$\mu_1 = \sqrt{\frac{m_u}{m_c}}, \quad \mu_2 = \sqrt{\frac{m_c}{m_t}}, \quad \nu_1 = \sqrt{\frac{m_d}{m_s}}, \quad \nu_2 = \sqrt{\frac{m_s}{m_b}}. \tag{21}$$

Then the mixing matrix (18) takes the following form

$$\begin{pmatrix} 1 & -\nu_1 + \mu_1 e^{i\sigma} & \mu_1 (\nu_2 e^{i\sigma} - \mu_2 e^{i\tau}) \\ -\mu_1 + \nu_1 e^{i\sigma} & \mu_1 \nu_1 + \mu_2 \nu_2 e^{i\sigma} + e^{i\tau} & \nu_1 (\mu_2 e^{i\sigma} - \nu_2 e^{i\tau}) \\ \nu_1 (\mu_2 e^{i\sigma} - \nu_2 e^{i\tau}) & \mu_2 \nu_1 + \mu_2 \nu_2 e^{i\sigma} + e^{i\tau} & \nu_2 e^{i\sigma} - \mu_2 e^{i\tau} \end{pmatrix}. \tag{22}$$

A detailed derivation of these relations are given in Ref[21]. A cancelation among two terms in the expression for $V_{cb}$ in (22) is needed. To achieve this one can choose

$$\sigma \sim \tau \sim -\frac{\pi}{2} \text{ this gives } V_{cb} \sim \sqrt{\frac{m_s}{m_b}} - \sqrt{\frac{m_s}{m_t}}. \tag{23}$$

It is easy to check that (23) makes the top quark mass too light to be experimentally true. Thus, it is worth asking the question whether if the Fritzsch texture were valid at the GUT scale instead, in other words, if the flavor
symmetries were exact only above the GUT scale, could a miracle of renormalization group evolution of the masses and mixing angles make the Fritzsch relations valid at low energy [22]. Here we study their idea.

The renormalization of the full $3 \times 3$ complex Yukawa matrices and their renormalization up to the GUT scale $M_X = 10^{16}$ GeV is what follows. Let us set our notations of mixing angles, the non-removable phase and eigenvalues of the Yukawa matrices. We adopt the parameterization [23]

$$
\begin{pmatrix}
  s_1 s_2 c_3 + c_1 c_2 e^{i \phi} & c_1 s_2 c_3 - s_1 c_2 e^{i \phi} & s_2 s_3 \\
  s_1 c_2 c_3 - c_1 s_2 e^{i \phi} & c_1 c_2 c_3 + s_1 s_2 e^{i \phi} & c_2 s_3 \\
  -s_1 s_3 & -c_1 s_3 & c_3
\end{pmatrix}.
$$

There is a detailed proof in Ref. [23] that in this parameterization eigenvalues $y_i$ of the Yukawa textures, three CKM mixing angles and the CP violating phase $\phi$ satisfy the renormalization group equations

$$
16\pi^2 \frac{d}{dt} \phi = 0,
$$
$$
16\pi^2 \frac{d}{dt} \ln \tan \theta_1 = -y_t^2 \sin^2 \theta_3,
$$
$$
16\pi^2 \frac{d}{dt} \ln \tan \theta_2 = -y_b^2 \sin^2 \theta_3,
$$
$$
16\pi^2 \frac{d}{dt} \ln \tan \theta_3 = -y_t^2 - y_b^2,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_u = -c_t^u g_t^2 + c_b^u g_b^2 + y_u^2 \cos^2 \theta_2 \sin^2 \theta_3,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_c = -c_t^c g_t^2 + c_b^c g_b^2 + y_c^2 \sin^2 \theta_2 \sin^2 \theta_3,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_t = -c_t^t g_t^2 + 6y_t^2 + y_t^2 \cos^2 \theta_3,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_d = -c_t^d g_t^2 + y_d^2 \sin^2 \theta_1 \sin^2 \theta_3 + 3y_b^2 + y_r^2,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_s = -c_t^s g_t^2 + y_s^2 \cos^2 \theta_1 \sin^2 \theta_3 + 3y_b^2 + y_r^2,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_b = -c_t^b g_t^2 + y_b^2 \cos^2 \theta_3 + 6y_b^2 + y_r^2,
$$
$$
16\pi^2 \frac{d}{dt} \ln y_e = -c_t^e g_t^2 + 3y_e^2 + y_r^2.
$$
\begin{align*}
16\pi^2 \frac{d}{dt} \ln y_\mu &= -c_i g_i^2 + 3y_\mu^2 + y_\tau^2 , \\
16\pi^2 \frac{d}{dt} \ln y_\tau &= -c_i g_i^2 + 3y_\mu^2 + 4y_\tau^2 . \tag{25}
\end{align*}

Multipliers $c_i$ of gauge couplings in (18) are well known. They can be found in [23]. We have solved these one-loop equations numerically using Mathematica NDSolve subroutine. The flow chart follows this line. Taking all experimentally possible values of the eigenvalues but only central values of the angles at low energy, we have evolved the set to the GUT scale using (18). At the GUT scale predictions for $V_{cb}$ $V_{us}$ and $V_{ub}$ are calculated assuming that Fritzsch relations are valid only at the GUT scale and beyond. While translating predictions of CKM entries back to low energy using (18), we have used exact values of the angles not central values. Thereafter we have checked whether each individual value of masses and mixings remain within experimentally allowed ranges. For the strange quark mass values quoted in table 1 are used. For all other masses and mixings the experimental values are taken from the review of particle physics[14]. Our results are given in table 2

| $\alpha_s$ | $\tan \beta$ | $m_s(2 \text{ GeV})$ |
|-----------|--------------|-------------------|
| 0.118     | 2            | 59.90 MeV         |
| 0.118     | 10           | 61.52 MeV         |
| 0.118     | 20           | 63.05 MeV         |
| 0.118     | 30           | 66.90 MeV         |

Table 2: Our results are quoted for $m_t^{pole} = 173$ GeV. All other masses and mixings remain within the ranges quoted by the Review of Particle Physics.

In conclusion, implications of results of $n_f = 2$ unquenched lattice simulations of the strange quark mass in the context of the Fritzsch texture are studied. Previous calculations in this line exist in the literature. We have two new aspects. Because the combined effect of charge and color breaking and perturbative unitarity of the Yukawa couplings may rule out the large $\tan \beta$ scenario we have studied the low $\tan \beta$ scenario. Moreover, we have included corrections of supersymmetric origin in the study.
Supersymmetric corrections to the bottom quark mass and the low \( \tan \beta \) scenario goes hand in hand. This is in the sense that in the low \( \tan \beta \) regime Fritzsch texture demands a large Yukawa contribution to the bottom quark mass. This is partially canceled by the supersymmetric loop corrections. Thus, the original Fritzsch texture is consistent with experimental data if it holds at the GUT scale. The strange quark mass emerges in the range 60-70 MeV at \( \mu = 2 \text{ GeV} \) for the central values of \( \alpha_s = 0.118 \) and \( m_t^{pole} = 173 \text{ GeV} \). This range is consistent with \( n_f = 2 \) sea quark effect improved (unquenched) lattice simulations of the strange quark mass.

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