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PHENOMENOLOGY OF SINGLET QUARKS*

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ABSTRACT

As possible extensions to the Standard Model, singlet quarks have interesting and testable consequences. We collect and update the constraints from present data on their masses and mixings with conventional quarks. The CP asymmetries in B^0 decays can differ dramatically from Standard Model expectations. The d-type singlets are accommodated economically in grand unification scenarios as pieces of additional $\mathbf{5}+\mathbf{5}^*$ multiplets of $SU(5)$. The u-type singlets could arise in $\mathbf{10}+\mathbf{10}^*$ multiplets of $SU(5)$.

1. Introduction

Singlet quarks are color-triplet fermions whose left and right chiral components are both singlets with respect to the SU(2) weak isospin gauge group. As such they offer an interesting example of physics beyond the Standard Model (SM). They can mix with the ordinary quarks of the SM, and thereby impact a wide variety of experimental measurements, as they generate tree-level FCNC's, introduce unitarity violation in the SM CKM matrix, influence neutral meson-antimeson oscillations, and modify CP asymmetries. These objects can be produced by strong, electromagnetic and weak-neutral-current interactions, and produce interesting decay signatures[1]-[12]. They might be indirectly detected by looking for additional sources of flavor-changing neutral currents (FCNC), flavor-diagonal neutral currents (FDNC), and CP violation[3],[13]-[24]. We summarize here some important implications and constraints on the phenomenology of singlet quarks; an expanded version can be found in Ref. [25].

Consider first the case of one charge $-\frac{1}{3}$ singlet field, mixing with the three SM fields of this charge. The mass eigenstates d,s,b,x are then linear combinations of the three orthonormal linear combinations d'_L, s'_L, b'_L of left chiral components

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that are $\mathrm{SU}(2)_L$ doublet partners of the known $Q = \frac{2}{3}$ fields u_L, c_L, t_L; the remaining orthonormal combination x'_L is an $\mathrm{SU}(2)_L$ singlet. Then

$$\begin{pmatrix} d'_L \\ s'_L \\ b'_L \\ x'_L \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} & V_{ux} \\ V_{cd} & V_{cs} & V_{cb} & V_{cx} \\ V_{td} & V_{ts} & V_{tb} & V_{tx} \\ V_{od} & V_{os} & V_{ob} & V_{ox} \end{pmatrix} \begin{pmatrix} d_L \\ s_L \\ b_L \\ x_L \end{pmatrix} . \tag{1}$$

The transformation matrix V is then the generalization of the usual 3×3 Cabibbo-Kobayashi-Maskawa (CKM) matrix V_{CKM}. The SM unitarity constraints on V_{CKM} no longer apply. One extra singlet quark added to the usual three generations of the SM gives rise to a 4×4 matrix V to which the unitarity conditions $VV^\dagger = I$ and $V^\dagger V = I$ apply. SM unitarity constraints on the 3×3 CKM matrix give linear three-term relations that can be expressed graphically as triangle relations in the complex plane as shown in Fig. 1. With 4×4 mixing, they become four-term relations; e.g. for one $Q = -\frac{1}{3}$ singlet, we have

$$V^*_{ui} V_{uj} + V^*_{ci} V_{cj} + V^*_{ti} V_{tj} + V^*_{oi} V_{oj} = \delta_{ij}, \tag{2}$$

or again,

$$V^*_{id} V_{jd} + V^*_{is} V_{js} + V^*_{ib} V_{jb} + V^*_{ix} V_{jx} = \delta_{ij}. \tag{3}$$

For $i \neq j$ these are expressible as quadrangle conditions in the complex plane. The first three terms in each case, however, are precisely the three sides of a triangle if CKM unitarity holds (the most discussed example is Eq.(2) with $i = b, j = d$). Thus 4×4 unitarity replaces the CKM triangle relations by quadrangle relations. In Eq.(2) the fourth side of the quadrangle is $V^*_{oi} V_{oj} = -z_{ij}$, the FCNC coefficient[20]. In Eq.(3) the fourth side is $V^*_{ix} V_{jx}$, that occurs in certain flavor-changing box diagrams (see below).

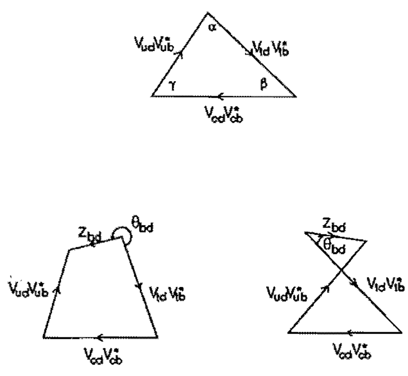

Fig. 1. Unitarity of the 3×3 CKM matrix implies triangle relations like the one shown. In 4×4 mixing cases the triangle relations become quadrangle relations.

It is therefore convenient to define the deviation from unitarity of the 3×3 CKM matrix by defining the quantities

$$z_{ij} \quad = \quad V^*_{ui} V_{uj} + V^*_{ci} V_{cj} + V^*_{ti} V_{tj} = \delta_{ij} - V^*_{oi} V_{oj} \; . \tag{4}$$

Similarly one can consider adding an u-type singlet to the SM quarks. Then one defines analogously the transformation matrix $\hat{V}$

$$\left(\bar{u}'_L \quad \bar{c}'_L \quad \bar{t}'_L \quad \bar{x}'_L \right) = \left(\bar{u}_L \quad \bar{c}_L \quad \bar{t}_L \quad \bar{x}_L \right) \begin{pmatrix} \hat{V}_{ud} & \hat{V}_{us} & \hat{V}_{ub} & \hat{V}_{uo} \\ \hat{V}_{cd} & \hat{V}_{cs} & \hat{V}_{cb} & \hat{V}_{co} \\ \hat{V}_{td} & \hat{V}_{ts} & \hat{V}_{tb} & \hat{V}_{to} \\ \hat{V}_{xd} & \hat{V}_{xs} & \hat{V}_{xb} & \hat{V}_{xo} \end{pmatrix} . \tag{5}$$

2. Constraints on singlet quark mixing

We give in this section the constraints on the singlet quark mixing matrices V and $\hat{V}$. We first define the source of the constraint, and summarize the resulting constraints at the end of this section.

Without imposing unitarity constraints, the CKM matrix elements lie in the following ranges [26]:

$$|V| = \begin{pmatrix} 0.9728 - 0.9757 & 0.218 - 0.224 & 0.002 - 0.005 & .. \\ 0.180 - 0.228 & 0.800 - 0.975 & 0.032 - 0.048 & .. \\ 0.0 - 0.013 & 0.0 - 0.56 & 0.0 - 0.9995 & .. \\ .. & .. & .. & .. \end{pmatrix} \tag{6}$$

The constraint due to the unitarity of the 4×4 matrix in Eq. (4) together with the bounds in Eq. (6) give rise to constraints on the individual entries in the matrix.

Singlet quarks provide additional contributions to meson-antimeson oscillations in two ways as shown in Fig. 2 for the B_d^0-$\bar{B}_d^0$ case: (1) the tree-level FCNC Z-exchange and (2) an additional contribution to the box graphs. The singlet quarks contribute in the two cases according to their charge as shown in Table 1.

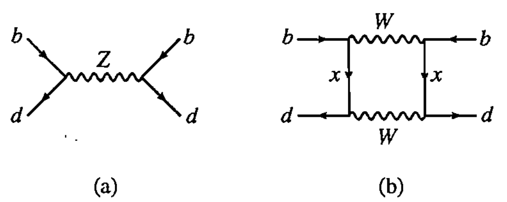

Fig. 2. Singlet quark mixing can give meson-antimeson oscillations via induced FCNC tree diagrams (a) and via box diagrams (b), illustrated here for the $B_d^0 - \overline{B}_d^0$ case.

The contributions from the Z-exhange graphs gives contraints on the z_{ij} as

Table 1: FCNC effects of singlet quarks

	$Q_x = -1/3$	$Q_x = 2/3$
K^0-$\bar{K}_0$ osc.	Z-exchange	box
D^0-$\bar{D}^0$ osc.	box	Z-exchange
B^0_d-$\bar{B}^0_d$ osc.	Z-exchange	box
B^0_s-$\bar{B}^0_s$ osc.	Z-exchange	box

they give rise to FCNCs at the tree-level through the coupling

$$\mathcal{L}_{\rm FCNC} \;=\; 1/2 g_Z \sum_{i\neq j} z_{ij}\bar{q}_{iL}\gamma^\mu Z_\mu q_{jL} \;, \;. \tag{7}$$

For the case of B^0_d-$\bar{B}^0_d$ oscillations, one obtains a contribution

$$|\delta m| = \frac{\sqrt{2}G_F m_B f_B^2 B_B \eta_B}{3}|z^2_{db}| \;, \tag{8}$$

where f_B is the B_d decay constant, B_B is the bag factor ($B=1$ is the vacuum saturation approximation) and $\eta_B \approx 0.55$ is a QCD factor.From the measurements [27] $|\delta m|_K = (3.51\pm0.02)\times10^{-12}\,$MeV, $|\delta m|_D < 1.3\times10^{-10}\,$MeV, $|\delta m|_{B_d} = (3.4\pm0.4)\times10^{-10}\,$MeV, one obtains bounds given at the end of this section.

The box-graph in the SM gives

$$|\delta m|_{\rm SM} = \frac{G_F^2 B f_B^2 m_B \eta_B}{6\pi^2}\,|V_{td}V^*_{tb}|^2_{\rm CKM}\,|I_B| \;, \tag{9}$$

where $I_B(x_t = m_t^2/M_W^2)$ is the box-integral (see e.g. Ref. [28])

$$I_B(x_t) = \frac{1}{4}M_W^2\left[x_t\left(1+\frac{9}{1-x_t}-\frac{6}{(1-x_t)^2}\right)-\frac{6x_t^3}{(1-x_t)^3}\ln x_t\right] \;. \tag{10}$$

I_B is a slowly varying function of its argument. Adding an extra singlet in the $x-t$ degenerate limit is accomplished by making the replacement

$$\left|\hat{V}_{xd}\hat{V}^*_{xb} + \hat{V}_{td}\hat{V}^*_{tb}\right|^2 = |V_{ud}V^*_{ub} + V_{cd}V^*_{cb}|^2_{\rm CKM} \tag{11}$$

using unitarity. However, $|V_{td}V^*_{tb}|_{CKM} = |V_{ud}V^*_{ub} + V_{cd}V^*_{cb}|_{\rm CKM}$, so the prediction for $|\delta m|$ is effectively unchanged in this x,t mass-degenerate limit. Since I_B is a slowly varying function of its argument, significant changes can arise only if x is much heavier than t. Similar conclusions apply to K^0-$\bar{K}^0$ oscillations. Additional constraints arise in D^0-$\bar{D}^0$ mixing and in ϵ_K; we refer the reader to Ref. [25] for details.

The tree-level FCNC Z couplings give rise to enhancements in some rare meson decays. Important constraints arise from the processes $K_L, B^0, D^0 \to \mu^+\mu^-$ and $B, D \to X\ell^+\ell^-$. On the other hand, deviations from the SM for $B\to s(d)\gamma$ decays are generally small[22, 25, 29].

The mixing of singlet quarks with the ordinary quarks of the SM changes the flavor-diagonal neutral current (FDNC) couplings. For the d-type singlet quark one obtains

$$\mathcal{L}_{\rm FDNC} = g_Z \sum_{i=d,s,b,x} \bar{q}_i\gamma^\mu Z_\mu\left[\frac{1}{4}z_{ii}(1-\gamma_5) - \frac{1}{3}\sin^2\theta_W\right]q_i \;. \tag{12}$$

Thus for the standard d, s, b quarks, mixing with x reduces direct left-handed FDNC by a factor $(z_{ii} - \frac{2}{3}\sin^2\theta_W)/(1 - \frac{2}{3}\sin^2\theta_W)$ and leaves right-handed FDNC unchanged. [Incidentally, since $b - x$ mixing reduces the $Z \to b\bar{b}$ coupling, it would make the discrepancy between the LEP data [30] and the SM prediction for the ratio worse[16, 18, 19]].

The global FDNC constraints are comprehensive enough to have useful repercussions via unitarity, for d-type singlet mixing. A comparison of all FDNC effects with the latest LEP and SLC data leads to global FDNC constraints (see final paper of Ref.[19]) listed below.

We summarize below the available bounds[25] on singlet quark mixing.

$Q = -\frac{1}{3}$ case		limit	origin
$\lvert V_{od}\rvert$	$\lesssim$	0.048	global FDNC
$\lvert V_{os}\rvert$	$\lesssim$	0.060	global FDNC
$\lvert V_{ob}\rvert$	$\lesssim$	0.045	global FDNC
$\lvert V_{ux}\rvert$	$\lesssim$	0.08	CKM + unitarity
$\lvert V_{cx}\rvert$	$\lesssim$	0.09	FDNC + unitarity
$\lvert V_{tx}\rvert$	$\lesssim$	0.09	FDNC + unitarity
$\lvert V_{ox}\rvert$	$\gtrsim$	0.996	FDNC + unitarity
$\lvert V_{os}\rvert\lvert V_{od}\rvert$	$\lesssim$	3×10^{-4}	$\epsilon, \delta m_K$(tree)
$\lvert V_{ob}\rvert\lvert V_{od}\rvert$	$\lesssim$	8×10^{-4}	δm_B(tree)
$\lvert V_{ob}\rvert\lvert V_{os}\rvert$	$\lesssim$	2×10^{-3}	$B \to \ell^+\ell^- X$
$\lvert V_{cx}\rvert\lvert V_{ux}\rvert$	$\lesssim$	$(1.3\mathrm{GeV})/m_x$	δm_D(box)
$\lvert Re(V_{od}^* V_{os})\rvert\lvert Im(V_{od}^* V_{os})\rvert$	$\lesssim$	3×10^{-10}	ϵ_K
$\lvert Re(V_{od}^* V_{os})\rvert$	$\lesssim$	7×10^{-6}	$K_L \to \mu\mu$

$Q = \frac{2}{3}$ case		limit	origin
$\lvert\hat{V}_{uo}\rvert$	$\lesssim$	0.049	global FDNC
$\lvert\hat{V}_{co}\rvert$	$\lesssim$	0.065	global FDNC
$\lvert\hat{V}_{to}\rvert$	$\lesssim$	1.0	unitarity
$\lvert\hat{V}_{xd}\rvert$	$\lesssim$	0.15	CKM + unitarity
$\lvert\hat{V}_{xs}\rvert$	$\lesssim$	0.56	CKM + unitarity
$\lvert\hat{V}_{xb}\rvert$	$\lesssim$	1.0	unitarity
$\lvert\hat{V}_{co}\rvert\lvert\hat{V}_{uo}\rvert$	$\lesssim$	9×10^{-4}	δm_D(tree)
$\lvert\hat{V}_{xd}\rvert\lvert\hat{V}_{xb}\rvert$	$\lesssim$	0.03	CKM + unitarity

3. CP asymmetries

The amount of CP violation in the SM is measured by the size of the unitarity triangle in Fig. 1. How this CP violation shows up in decays is determined by the angles of the unitarity triangle(s), which appear as CP asymmetries in decays to CP eigenstates. The angles shown in Fig. 1 are defined as

$$\beta \equiv \arg\left(-\frac{V_{cd}V_{cb}^*}{V_{td}V_{tb}^*}\right), \qquad \alpha \equiv \arg\left(-\frac{V_{td}V_{tb}^*}{V_{ud}V_{ub}^*}\right), \tag{13}$$

are directly measurable in B_d decays with $b \to c$ and $b \to u$ respectively. The prototype processes for measuring β and α in the SM are $B_d \to \psi K_S$ and $B_d \to \pi^+\pi^-$

respectively. In the presence of mixing with singlet quarks, the CP-asymmetries are no longer simply related to these angles, since there is an additional $B_d - \bar{B}_d$ oscillation contribution from tree-level Z-mediated graphs. Present information on the third generation couplings does not tell us much about the asymmetries. Future improved measurements of the CKM mixing angles at a B-factory, for example, will pin down the SM prediction more precisely[32].

The time-dependent CP asymmetry in the decay of a B_d^0 or $\overline{B}_d^0$ into some final CP eigenstate f is (assuming as usual $\Gamma_{12} << M_{12}$)

$$\frac{\Gamma(B_d^0(t) \to f) - \Gamma(\overline{B}_d^0(t) \to f)}{\Gamma(B_d^0(t) \to f) + \Gamma(\overline{B}_d^0(t) \to f)} = -\mathrm{Im}\ \lambda(B_d \to f)\ \sin(\delta m\ t)\ , \tag{14}$$

where δm is the (positive) difference in meson masses, the mesons states evolve from flavor eigenstates B_d^0 and $\overline{B}_d^0$ at a time $t = 0$, and $\mathrm{Im}\ \lambda(B_d \to f)$ is the time-independent asymmetry.

The allowed range[31] for the CP-asymmetry in the SM for the quantity $\mathrm{Im}\ \lambda(B_d \to \psi K_S)$ is shown in Fig. 3 as a shaded band (for fixed m_t). The expectations for the same quantity in the presence of a d-type singlet quark is shown in Fig. 1 for different values of the parameters[22]

$$\delta_d \equiv \left|\frac{z_{bd}}{V_{td}V_{tb}^*}\right|\ , \qquad \theta_{bd} \equiv \arg\left[\frac{z_{bd}}{V_{td}V_{tb}^*}\right]\ , \tag{15}$$

which characterize the unitarity violation. The effect of singlet quarks on CP asymmetries can be dramatic[20, 22, 32]. One notices that the CP asymmetry $\mathrm{Im}\ \lambda(B_d \to \psi K_S)$ is negative (with our conventions) in the SM, but with sufficiently large δ_d one can obtain positive values[20, 22].

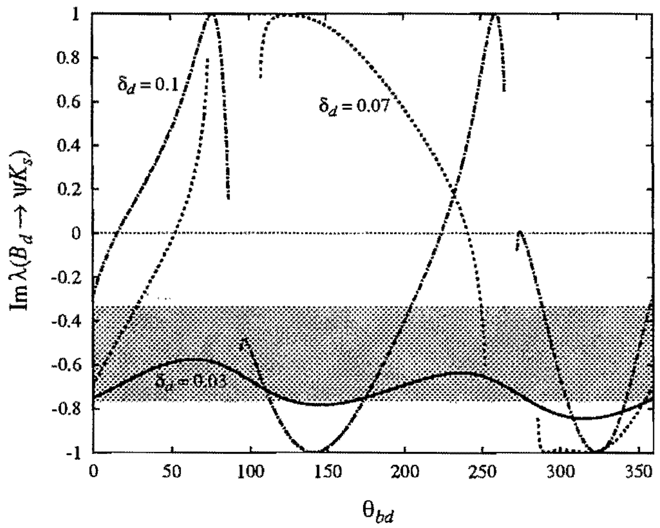

Fig. 3. The CP asymmetry $\mathrm{Im}\ \lambda(B_d \to \psi K_S)$ in the presence of down-type singlet quarks. The band indicates the present uncertainty in the SM prediction. For some values of δ_d and θ_{bd} (defined by Eq. (15)) there are no solutions as the unitarity quadrangle cannot be made to close.

The CP-asymmetry Im $\lambda(B_d \to \pi^+\pi^-)$ is not as well constrained in the SM; but given well-determined CKM elements, deviations from SM predictions could be significant and could provide evidence for singlet quarks[25, 32].

4. GUT sources of singlet quarks

Singlet quarks cannot simply be added by themselves to the SM or MSSM and gauge coupling unification still be successful. Considered alone they do not introduce gauge anomalies, but they change the running of the $SU(3)$ and $U(1)$ couplings but not the $SU(2)$ coupling. For down-type singlets, this can be remedied by adding more fermions to fill out the **5** and **5*** representations of $SU(5)$, or a **10** of $SO(10)$[33, 34]; see Fig. 4. This exotic matter together with the SM matter content fits neatly into the **27** representation of E_6[33, 35] (which decomposes to $\mathbf{10} + \mathbf{5}^* + \mathbf{1} + \mathbf{5} + \mathbf{5}^{*\prime} + \mathbf{1}$ in an $SU(5)$ subgroup). Adding extra complete multiplets of $SU(5)$ preserves (at the one-loop level) the successful unification of gauge couplings in the MSSM, since a complete multiplet contributes equally to the evolution of each coupling. However, more than three generations of exotic matter will destroy asymptotic freedom for α_3 at one-loop.

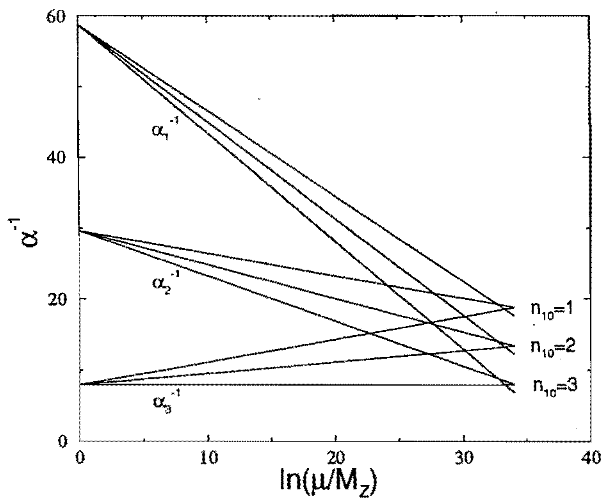

Fig. 4. One-loop gauge coupling evolution with the addition of different numbers of light **10** multiplets of $SO(10)$ to the MSSM. Successful gauge coupling unification is preserved with the addition of complete **10**-plets.

An up-type singlet quark is not contained as elegantly in GUT Models; it does not appear in the smallest representations, and its role is less clear. As a minimal prescription, it can be introduced by adding one extra light **10** and one **10*** representation of $SU(5)$ that get their mass from an $SU(5)$ singlet Higgs boson; this implies extra vector-doublet quarks and a vector-singlet charged lepton too, preserving MSSM gauge coupling unification with $b_3 = 0$ at one loop. Less minimally, it can also be realized in the $SO(10)$ group with an extra light **45** (adjoint) representation (which decomposes to $\mathbf{24} + \mathbf{10}^* + \mathbf{10} + \mathbf{1}$ in an $SU(5)$ subgroup), but this leads to nonperturbative gauge couplings at the GUT scale if the entire **45** is required to be light.

Fixed points play a role in these extended model, since the extra matter content implies larger gauge couplings. Hence the top quark and d-type singlet(s) masses could possibly be determined by the gauge couplings and the associated vevs. However, the Yukawa unification condition $\lambda_b(M_G) = \lambda_\tau(M_G)$ becomes harder to accomodate, and fails in the E_6 model with three light generations.

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