Superheavy Light Quarks and the Strong P, T Problem

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Abstract

New superstrong forces, analogous to QCD but featuring a larger mass scale, should they exist, offer new perspectives on the strong P, T problem. If the superstrong dynamics supports confinement without chiral symmetry breaking, then one might implement the “massless quark” solution in a phenomenologically acceptable way, by using a massless quark that is always confined within superheavy particles, and is therefore effectively superheavy: a cryptoquark. Assuming confinement and chiral symmetry breaking from the superstrong dynamics, we find a new mechanism to generate an axion field without introducing new fundamental scalar fields.

Our present-day core theories – general relativity, and the $SU(3) \times SU(2) \times U(1)$ gauge theories of strong, weak, and electromagnetic interactions – offer a profound explanation of the origin of (approximate) invariance of fundamental physics under the time reversal operation $T$. As emphasized by Kobayashi and Maskawa [1], (approximate) $T$ symmetry arises as an accidental consequence of more basic principles, i.e. the principles of relativistic quantum field theory, together with the observed multiplet structure of matter in $SU(3) \times SU(2) \times U(1)$. Since $T$ symmetry is emphatically not manifested in macroscopic observables, it might otherwise appear as a gratuitous and even problematic feature of our world-description.
At first sight the three-family version of $SU(3) \times SU(2) \times U(1)$ appears to allow only a single $T$-violating parameter. That parameter is a weak mixing angle involving transitions to the third generation. The effects it induces are predicted to be visible only in $K - \bar{K}$ mixing and in processes involving the third generation, and in fact they have been observed.

There is, however, a feature of that explanation which remains unsatisfactory. Deeper investigation reveals an additional possibility for $T$ violation, bringing in an independent parameter $[2]$. The additional parameter is the coefficient of a color gauge field version of $\Delta \mathcal{L} \propto \vec{E} \cdot \vec{B}$, which can appear as an interaction term in the theory’s Lagrangian density. Since that interaction is a total divergence, it does not appear in the classical equations of motion. But the quantity of which it is a divergence is not gauge invariant, so in the quantum theory it can (and does) exhibit singular fluctuations, leading to surface terms. Thus the $\vec{E} \cdot \vec{B}$ interaction term can (and does) influence the quantum theory.

One can parameterize the coefficient of the new interaction in terms of a variable $\theta$, which is periodic modulo $2\pi$. $\theta$ is odd under $T$ (and $P$), so those symmetries are broken unless

$$\theta \equiv 0, \pi \pmod{2\pi}$$

The new interaction can be calculated to induce an electric dipole moment for the neutron, which violates both $P$ and $T$ invariance. No such moment has been observed, despite very sensitive searches. That result, together with some special considerations that exclude $\theta \approx \pi$, lead to the bound $[3]$

$$|\theta| < 10^{-10} \pmod{2\pi}$$

The smallness of this number does not follow from general principles. It constitutes a remaining gap in our understanding, often called the strong $P,T$ problem.

One can visualize the role of $\theta$ by introducing an effective interaction, the 't Hooft vertex, depicted in Figure 1. The blob in the 't Hooft vertex indicates a gauge field configuration with unit topological charge, and the fermion lines emanating from it indicate fermion emission and absorption. One left-handed quark of each flavor is absorbed, and one right-handed quark of each flavor is emitted. (This representation is schematic, but adequate to our purposes.) The interaction contains an overall factor $e^{i\theta}$. Now if the $u$ quark were massless, then it would be possible to redefine $u_L$ by a phase factor without changing the form of any term in the Lagrangian other
than the ’t Hooft vertex. In particular, we could use that freedom to absorb the factor $e^{i\theta}$ into a re-definition of $u_L$. In the new variables we would have $\theta = 0$. Thus the standard $P$ and $T$ operations would become manifest symmetries of the effective interaction, and the strong $P,T$ problem would be solved.

![Diagram](image)

Figure 1: Gauge configurations with unit topological charge are weighted with the phase parameter $e^{i\theta}$. In the presence of such gauge fields charged fermion fields have zero modes whose quantum numbers, including their chiral structure, implement appropriate anomaly equations (Atiyah-Singer index theorem). The ellipsis indicates additional quark species, beyond $u$.

This proposed “zero mass” solution is almost as old as the strong $P,T$ problem itself. (See, for example, [8].) But it appears that in reality $m_q \neq 0$ for every one of the known quarks, implying that Nature has not chosen to use that solution [4].

It is logically possible, however, that there are massless (either literally, or in a sense to be defined below) “quarks” $Q$ – that is, spin $\frac{1}{2}$ fermions carrying non-trivial color quantum numbers – yet to be discovered. It will be convenient to have a name for such particles, and we shall call them cryptoquarks. Indeed, there is no necessary relation between the current mass of confined quarks and the mass of the lightest particles which contain them. If
there is a new strong interaction $G$, with dynamics resembling that of QCD’s $SU(3)$ but on a much larger mass scale, and massless spin $\frac{1}{2}$ fermions carrying both $G$ and $SU(3)$ quantum numbers, then direct evidence for the existence of those fermions (or, for that matter, of $G$) could be difficult to obtain. One must exceed the threshold for production of the (encrypting) particles containing the cryptoquarks, or the crossover, at higher energy, to liberation of $G$ jets. (Related ideas are expressed in \cite{5} \cite{6}.)

The mechanism whereby cryptoquarks solve the strong $P,T$ problem is essentially identical to what we described above, with $u \to Q$, apart from one complication. (See Figure \ref{fig:2}.) The complication is, that $Q_L, Q_R$ will also appear in a separate ’t Hooft vertex, associated with $G$ (and a phase angle $\theta'$), so it is not correct that we can redefine their relative phase, and thus alter $\theta$, without changing anything else in the Lagrangian. To alleviate that difficulty, we can introduce another massless spin-$\frac{1}{2}$ fermion $S$ with different quantum numbers from $Q$ under $SU(3)$ and/or $G$. $S$ will generally appear with different multiplicities from $Q$ in the two vertices. If the multiplicity pairs are linearly independent, then we can set both $\theta$ and $\theta'$ to zero by appropriate re-definitions of $Q_L$ and $S_L$. The simplest possibility is for $S$ to transform under $G$ but not under $SU(3)$.

The suggested model – an additional gauge sector $G$ with massless (vectorlike) fermions $Q$ and $S$, which are respectively in the fundamental and singlet representations of $SU_c(3)$ – possesses chiral symmetries. If the superstrong dynamics do not break these chiral symmetries, then all cryptohadrons will be heavy and presumably unobservable. However, it appears more likely that the confining dynamics of the $G$ field spontaneously breaks these chiral symmetries, giving rise to Nambu-Goldstone modes, which can infest low-energy physics. Therefore we should investigate the pattern of symmetries and symmetry breaking. For first orientation, let us consider $SU_c(3)$ gauge interactions as a perturbation. Then $G$ possesses four cryptoquarks; $S$ and the 3 color components of $Q$. At the classical level, it has a symmetry

$$SU_L(4) \times SU_R(4) \times U_V(1) \times U_A(1).$$

The $U_A(1)$ symmetry is anomalous, and should not be included amongst the symmetries. (See below.) Conventional (QCD like) confinement triggered by the $G$ gauge fields gives rise to spontaneous breaking according to

$$SU_L(4) \times SU_R(4) \times U_V(1) \to SU_V(4) \times U_V(1).$$

At this level, we obtain 15 associated Nambu-Goldstone (cryptopion) modes.
Figure 2: On the left, the 't Hooft vertex for color $SU(3)$, with the presence of $Q$ highlighted. The ellipsis indicates other quark species, possibly including additional samples of $Q$, and of $S$. On the right, the 't Hooft vertex for a hypothetical new superstrong interaction with gauge group $G$. Gauge configurations with unit topological charge are weighted with the phase parameter $e^{i\theta'}$. The ellipsis indicates additional quark species, possibly containing additional samples of $S$ and of $Q$.

Now let us restore the $SU_c(3)$ color gauge interactions. The initial symmetry is reduced, at the classical level, to an $SU_V(3)$ acting on $Q$, $U_V(1) \times U_V(1)$ of $Q$ and $S$ number, and an axial $U_A(1)$ which acts on both $Q$ and $S$, in such a way as to avoid the $G$ anomaly. All rotations between $Q$ and $S$ fields are eliminated because of their different $SU_c(3)$ interactions. And $SU_c(3)$ also breaks the axial part of the $SU_L(3) \times SU_R(3)$ symmetry explicitly. Further, the residual $U_A(1)$ has a color $SU_c(3)$ anomaly. Plausibly there will still be, as a product of confining $G$ dynamics, an $SU_c(3)$ color octet of pseudoscalar bosons of type $Q_{LA} Q_{RA}^\beta$, and a similar $3 + \bar{3}$ involving $SQ$ and $QS$. They will be lighter than typical $G$ hadrons by factors of the kind $\alpha_s^p(\Lambda_G)$, where $\Lambda_G$ is the dynamical scale of $G$ and $p$ is a small positive number. Those masses could be very large by contemporary standards, and then the phenomenological impact of such massive particles would then be slight.

The residual $U_A(1)$ anomaly has a topological character, however, and
is associated with low-energy QCD dynamics. In a semiclassical framework, its physical effect arises solely from instantons. The would-be Nambu-Goldstone boson associated with breaking of that approximate symmetry plays the role of a QCD axion, with its dynamical scale set by the dynamical chiral symmetry breaking scale of $G$.

The mechanism thus exemplified is of course much more general than the particular model we chose to analyze. It rests on the spontaneous breaking by $G$ of a symmetry that is anomalous under color $SU_c(3)$, but not colored. The associated would-be Nambu-Goldstone fits the conventional phenomenological profile of a high-scale axion $[10,11]$, but its origin is significantly different, and perhaps more attractive. It does not require the introduction of any fundamental scalar field, and avoids associated mass hierarchy issues.

Our example possesses two vectorlike $U(1)$ symmetries, associated with cryptobaryon numbers. The associated cryptobaryons are absolutely stable, which in some cosmological scenarios could prove problematic. Simple variants of the model need not have this feature.

A noteworthy connection between massless cryptoquarks and more conventional axion schemes emerges if we expand the model just mentioned to include two additional complex scalar fields $\phi_Q, \phi_S$ with two Peccei-Quinn type symmetries $[7]$ of the type

$$
\begin{align*}
Q_L & \rightarrow e^{i\alpha} Q_L \\
S_L & \rightarrow e^{i\beta} S_L \\
\phi_Q & \rightarrow e^{i\alpha} \phi_Q \\
\phi_S & \rightarrow e^{i\beta} \phi_S
\end{align*}
$$

These symmetries support Yukawa couplings of the form

$$
\Delta\mathcal{L} = g_Q\phi_Q Q_L Q_R + g_S\phi_S S_L S_R + \text{h.c.}
$$

which also violate global axial $Q$ color explicitly.

If neither $\phi_Q$ nor $\phi_S$ were to develop a vacuum expectation value, we would in essence recover the cryptoquark model just discussed, augmented with some gratuitous scalar fields. If one but not the other develops a vacuum expectation value, we shall still have the cryptoquark mechanism for $G$, in addition to a more-or-less conventional QCD axion $[8,9,10,11]$. If both develop vacuum expectation values, we shall have two axion-like particles, one more-or-less conventional, the other putting the large energy scale $\Lambda_G$ into play (in place of $\Lambda_{\text{QCD}}$), and hence with a larger mass than
might be expected, given its coupling strength to color gluons. (It might, for example, weigh hundreds of GeV and be highly unstable.) That last case is favored, since the plausible formation of $\langle \bar{Q}_L Q_R \rangle$ and $\langle \bar{S}_L S_R \rangle$ condensates introduces linear terms in the potentials for $\phi_Q, \phi_S$, and induces non-zero vacuum expectation values for them.

The basic mechanisms for addressing the strong $P,T$ problem here described can be implemented in many ways. In particular, they are not inconsistent with promising ideas about unification of quantum numbers and couplings, nor with low-energy supersymmetry. Were a new superstrong force to be discovered, it would be interesting to check whether these mechanisms are operative. Conversely, one might be driven, from considerations within this circle of ideas, to infer the existence of a new superstrong force.

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