Method for Constructing Models with Strong *CP* Invariance

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A method is given which leads to a wide class of models with strong *CP* invariance. These models contain a spontaneously broken U(1) symmetry. If this symmetry were global, its breaking would lead to a massless scalar. Because the U(1) is anomaly free, it can be gauged, leading by the Higgs mechanism to a massive gauge boson, the aspon, which might be detectable.

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Quantum chromodynamics (QCD), the theory of strong interactions in particle and nuclear physics, experienced[1,2] both a major advance and a significant setback in 1976. The advance came when it was realized that the long-standing problem of the η meson mass arises from topological field configurations or instantons and hence resolved the so-called U(1) problem. At the same time, the discovery of instantons created the strong *CP* problem which remains unresolved.

The strong *CP* problem is that instantons introduce a new parameter $\theta_{\rm QCD}$ into QCD and that this parameter must be fine tuned to 1 part in 10^{10} to avoid disagreement with experiment. In particular, instanton effects may be summarized by an additional term in the QCD Lagrangian, with coefficient $\theta_{\rm QCD}$, which violates *P* and *CP* conservation. Experimental strong *CP*-violation effects depend on $\bar{\theta}=\theta_{\rm QCD}+\theta_{\rm QFD}$, where $\theta_{\rm QFD}$ is the phase of the determinant of the quark mass matrix.

The present Letter discusses the strong *CP* problem from a new perspective, provides a new solution, and should facilitate appropriate model building.

We first recall two approaches which have attracted considerable attention previously. The first,[3] and the more popular, is to introduce a color-anomalous $\mathrm{U(1)_{PQ}}$ which allows $\bar{\theta}$ to relax to zero; the spontaneous breaking of this $\mathrm{U(1)_{PQ}}$ gives rise to a light boson, the axion, which obtains a mass from instanton effects. A second approach is to assume *CP* symmetry of the Lagrangian so that $\theta_{\rm QCD}=0$; after spontaneous breaking of *CP*, the value of $\theta_{\rm QFD}$ is kept small by arranging a real determinant of the quark mass matrix at lowest order.[4]

Both methods[3,4] appear to offer acceptable solutions. One reason that more attention has been given to the axion scenario is that it gives rise to so many additional questions which require further research. A second reason is that there is a systematic method to construct Peccei-Quinn models. As searches for a physical axion remain frustrated, it is worth examining what other observable phenomena or particles might be associated with solution of the strong *CP* problem.

In order to set the scene, and to introduce a more general approach to model building, let us consider one family of quarks and their (T_3,Y) values under the electroweak group:

$$(-\tfrac{1}{2},\tfrac{1}{6})\ d_L,\quad (0,\tfrac{1}{3})\ \bar{d}_L\,,$$
$$(\tfrac{1}{2},\tfrac{1}{6})\ u_L,\quad (0,-\tfrac{2}{3})\ \bar{u}_L\,.$$

We introduce in our model a $\mathrm{U(1)_{new}}$ symmetry and assign charge $Q_{\rm new}=0$ to all of the above quark states and to the leptons, although the latter do not play a significant role in solving strong *CP*. The second and third families have the parallel assignments under the same $\mathrm{U(1)_{new}}$.

In our model there is also a real representation of exotic "heavy" quarks corresponding to a complex representation C and its conjugate $\bar{C}$. In $\bar{C}$ the exotic heavy quarks have quantum numbers exactly like some of the usual quarks; for example, in $\bar{C}$ there may be one doublet

$$(-\tfrac{1}{2},\tfrac{1}{6})\ D_L\,,$$
$$(\tfrac{1}{2},\tfrac{1}{6})\ U_L\,.$$

These have charge $Q_{\rm new}=+h$. In representation C we shall then have

$$(\tfrac{1}{2},-\tfrac{1}{6})\ D_L^C\,,$$
$$(-\tfrac{1}{2},-\tfrac{1}{6})\ U_L^C\,.$$

These have $Q_{\rm new}=-h$.

The Higgs sector has one complex doublet

$$\phi\ \ (+\tfrac{1}{2},-\tfrac{1}{2}),\quad Q_{\rm new}=0\,,$$

and two complex singlets

$$\chi_{1,2}\ \ (0,0),\quad Q_{\rm new}=+h\,.$$

The gauge group is $\mathrm{SU(3)_3\times SU(2)_{\mathit{L}}\times U(1)_{\mathit{Y}}}$ $[\times\mathrm{U(1)_{new}}$ if it is gauged]. In breaking the symmetry, we give a real vacuum expectation value (VEV) to ϕ and *complex* VEVs to $\chi_{1,2}$ with a nonvanishing relative phase.

The Lagrangian contains bare mass terms $M(U_L^c U_L + D_L^c D_L)$ for the extra quarks. The allowed Yukawa couplings include $\bar{u}_L^i u_L^j \phi$, $\bar{d}_L^i d_L^j \phi$ for the families and $U_L^c u_L^i \chi_\alpha$, $D_L^c d_L^i \chi_\alpha$ coupling light quarks to C heavy quarks ($\alpha=1,2$ and $i=1,2,3$). Because the families have no Yukawa couplings to $\bar{C}$ exotics, the quark mass matrix determinant arising from spontaneous symmetry breaking is real at lowest order; it has the required texture.[4]

We do not allow terms in the Higgs potential which explicitly break $\mathrm{U(1)_{new}}$. Disallowed terms include $\bar{\phi}\phi\chi$, $\bar{\phi}\phi\chi^2$, $\chi^2\chi^3$, and χ^4. If any of these terms are present, $\mathrm{U(1)_{new}}$ is explicitly broken and the model can have $\bar{\theta}=0$ at tree level only in the very special cases[4] where, e.g., we choose particular representations of a grand unified group such that the quark mass matrix is real.

Without explicit breaking of the $\mathrm{U(1)_{new}}$, there is correct texture at tree level; the mass matrix has the tree-level texture (F = family)

$$(F\ \bar{C}\ C)\begin{pmatrix} \text{real} & 0 & \text{complex} \\ \text{complex} & \text{real} & 0 \\ 0 & 0 & \text{real} \end{pmatrix}\begin{pmatrix} F \\ C \\ \bar{C} \end{pmatrix}.$$

Thus $\theta_{\mathrm{QFD}}=0$ at tree level. If we assume CP symmetry of the Lagrangian then $\theta_{\mathrm{QCD}}=0$ also. In this case $\bar{\theta}=0$ at tree level and will be nonzero by a small amount through radiative corrections; this is consistent with experiment if the off-diagonal F-C-χ Yukawa couplings are $<10^{-4}$.

If we assume that it is merely a global symmetry, the spontaneous breaking of $\mathrm{U(1)_{new}}$ gives rise to a massless boson, the χ. While an axion couples to a color-anomalous current, the χ couples to a nonanomalous current. Otherwise, their origins are similar: Both arise from spontaneously breaking a chiral U(1) which is imposed to solve the strong CP problem. The χ couples to a nonanomalous current and hence does not gain mass from instantons.

The χ resembles, in some respects, the familon[5] of Reiss and Wilczek. As a Goldstone boson, the χ couples derivatively and hence its static potential falls as R^{-4} with distance like the familon potential but with an additional suppression factor since χ couple to normal matter off diagonally. Between two protons, the χ force is stronger than gravity only for $R<(10^5/F)$ cm, where F (in GeV) is the scale characterizing the χ. If $F<10^5$ GeV, it should be possible to detect the χ force directly.

The spontaneously broken $\mathrm{U(1)_{new}}$ gives rise to global cosmic strings which might be astrophysically significant in processes such as galaxy formation.[6] The VEV of χ actually breaks a $\mathrm{U(1)_{new}}\times Z_2$ symmetry, where Z_2 is associated with the CP transformation. Breaking Z_2 gives rise to cosmological domain walls which must be removed by inflation since such walls would conflict with the present energy density of the Universe. The energy scale, or temperature, at which inflation takes place is limited only by the requirement that baryogenesis must occur *after* inflation. Since baryogenesis is now believed to be possible at the electroweak scale,[7] we may take $|\langle\chi\rangle|$ to be as low as only a few TeV. We then assume inflation takes place not far above the weak scale.[8] To summarize, the cosmological scenario is the following: (i) break $\mathrm{U(1)_{new}}$ and CP at roughly 2 TeV (see below for the justification of this choice), (ii) inflate below 2 TeV but above the weak scale (using perhaps an extended inflationary scheme involving one of the new scalar singlets), and (iii) generate baryon number at the weak scale via sphalerons.

One terrestrial experiment where the χ might be detected is the search for rare decays such as $\mu\to e\chi$, $K^+\to\pi^+\chi$. Here the situation is again similar but not identical to the case of familons;[5] the branching ratios for χ decay are smaller because a heavy exotic fermion is involved. and they are therefore consistent with experimental bounds.

There is a much more attractive possibility than a global $\mathrm{U(1)_{new}}$. Because it is anomaly free, we may *gauge* $\mathrm{U(1)_{new}}$ and hence avoid the χ boson; instead, there will be a massive gauge boson A^0, the aspon, which couples directly only to exotic fermions and indirectly via nondiagonal mixing to quarks and leptons. The gauged charge is, in this case, the exotic quark number which is $+h$ for exotic quarks and $-h$ for exotic antiquarks. Observability of the aspon will depend on its mass and the gauge coupling constant (g) which are related by $m(\text{aspon})\sim g\langle|\chi|\rangle$. The aspon could be quite light, and therefore interesting phenomenologically.

This model with a gauged $\mathrm{U(1)_{new}}$ we shall call the aspon model. It introduces several new parameters to the standard model. For the case of a heavy quark doublet, a detailed analysis of the consistency with phenomenology has been made.[9] The results of Ref. 9 may be summarized as follows. New flavor-changing neutral currents are naturally suppressed to levels below those already in the standard model. Keeping $\bar{\theta}$ sufficiently small and arranging to agree with the weak CP parameters $\varepsilon=2.258\times10^{-3}$ and $|\varepsilon'/\varepsilon|<4\times10^{-3}$ leads to the conclusion that CP must be spontaneously broken at a scale (characterized by the VEV of χ) below 2 TeV. The aspon mass is then less than 600 GeV if we assume $g_{\mathrm{new}}\le e$. The mass M of the heavy exotic quarks must be less than 530 GeV.

Compared to the most familiar solution of the strong CP problem,[3] the present model replaces an invisible axion by a visible aspon. Further, the fact that $\mathrm{U(1)_{new}}$ is local seems more appealing than the global symmetry of Ref. 3 since global symmetries are open to the well-known conceptual difficulty that they envisage a transformation made not only here and now but "behind the moon next week."[10]

It is interesting to tie together the solution of strong CP with the weak CP parameters of the Kobayashi-Maskawa matrix relevant to the neutral-kaon system and to B physics.

Although the aspon couples directly only to the exotic quarks, it does couple through off-diagonal mixing to the light quarks and if its mass is in an accessible range it should be observable; we look forward to its discovery.

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