Phenomenology of the Aspon Model of *CP* Violation

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(Received 26 September 1991)

The aspon model of spontaneous *CP* violation solves the strong *CP* problem and gives a new mechanism for weak *CP* violation. It predicts a vector quark doublet Q and a gauge boson A. The $Q\bar{Q}$ production at Fermilab and the Superconducting Super Collider (SSC) is calculated, production of a $Q\bar{Q}A$ final state at the SSC is computed, and signatures of Q and A are discussed.

PACS numbers: 12.15.Cc, 11.30.Er, 14.80.Er, 14.80.Pb

The standard model of elementary particle interactions is the most satisfying achievement of particle theory in the last three decades. It explains a great deal but contains twenty arbitrary parameters to be fitted phenomenologically. One feature of particle theory which is not satisfactorily explained in the standard model is *CP* violation in the neutral kaon decay. It can be *accommodated* within the Cabibbo-Kobayashi-Maskawa (CKM) matrix of flavor mixing between three families of quark doublets; indeed, in a way this *necessitates* three families. What is wrong with attributing *CP* violation to the CKM phase (δ) in the mixing matrix?

There is one other parameter in the standard model associated in an important manner with *CP* violation, namely, the angle $\bar{\theta}$ which must be extremely close to zero to avoid too large *CP* violation in the strong interactions. The CKM phase (δ) in the standard model offers no clue to why $\bar{\theta}$ is so small in magnitude.

In the aspon model [1,2] the mechanism of weak *CP* violation is by exchange of the heavy gauge boson A of $U(1)_{\rm new}$, a gauge group which is also able to solve the problem of strong *CP* conservation. In linking the two parameters δ and $\bar{\theta}$ of the standard model, this mechanism of *CP* violation is therefore more satisfactory.

In the model, a heavy vector quark doublet $Q=(U,D)$ has Dirac mass M and Yukawa couplings $h_i^{(a)}$ ($i=1,2,3$; $a=1,2$) with the three light left-handed quark doublets $q_{iL}=(u_i,d_i)_L$ and two singlet scalar fields $\chi^{(a)}$. When the $\chi^{(a)}$ develop complex vacuum values which spontaneously break *CP*, the resulting mass matrix depends on the mixing parameters $x_i=M^{-1}\sum_a h_i^{(a)}\langle\chi^{(a)}\rangle$.

In the present work we are assuming that these light-heavy quark mass mixing parameters x_i are sufficiently small, $x_i<0.1$, that a perturbative expansion in powers of x_i is rapidly convergent. Larger x_i may be allowed. (We thank E. Carlson for a very informative discussion on this point.) In this small-x_i perturbative regime, the aspon model predicts the existence of the nearly degenerate doublet $Q=(U,D)$ of color-triplet fermions (called *quirks*) with Dirac mass $M<300$ GeV, as well as the gauge boson A with $M_A<600$ GeV. The gauge boson corresponds to a gauge group $U(1)_{\rm new}$ under which all particles within the standard model are neutral and whose charge is carried by the quirk fields Q and by the singlet scalars $\chi^{(a)}$. The production and signatures of these hypothetical particles are the subject of the present Letter. It is a noteworthy feature of this perturbative aspon model that the scale of new particles is accessible to the Superconducting Super Collider (SSC). In other solutions to the strong *CP* problem, e.g., the Peccei-Quinn solution [3] or the Barr-Nelson [4] solution, the scale of new physics is the grand unification scale or nearly so, and not testable at the SSC.

We begin with the production of $Q\bar{Q}$ pairs at hadron collider pp or $p\bar{p}$. The relevant center-of-mass energies are $\sqrt{s}=2$ TeV (Fermilab) and 40 TeV (SSC). The production mechanism is closely parallel to that for $t\bar{t}$, where t is the top quark, and hence is dominated by gluon fusion. Adding the two different gluon fusion graphs and using the structure functions of Ref. [5] gives the results shown in Fig. 1 for $Q\bar{Q}$ production at $\sqrt{s}=40$ and 2 TeV. At Fermilab energy, the cross section falls below 1 pb for $M\geq 160$ GeV; as expected, the production capability is the same as that for the t quark. As we shall discuss

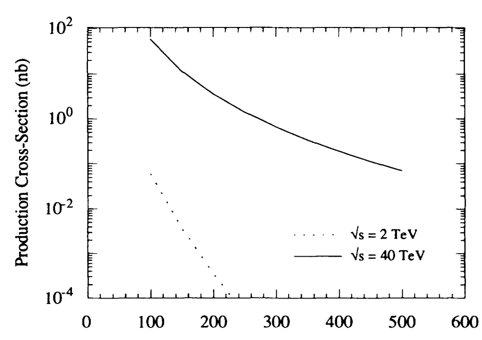

FIG. 1. Cross section for pp and $p\bar{p}\rightarrow Q\bar{Q}$ from gluon fusion at $\sqrt{s}=40$ and 2 TeV, respectively.

later, decay of the U quirk resembles that of the t quark so that the limit $m_t > 89$ GeV (Ref. [6]) may be taken immediately to imply $M > 89$ GeV also. At SSC, we see that even for M at its maximum value of 300 GeV, the cross section is almost 1 nb, which translates to 10^7 events/yr at the planned SSC luminosity of 10^{33} $\text{cm}^{-2}\text{s}^{-1}$.

The aspon particle is most copiously produced at the giant proton collider in the final state $Q\bar{Q}A$ (see Fig. 2). In Fig. 3 is shown the cross section in pp at $\sqrt{s}=40$ TeV for $M=100,200,300$ GeV. For typical values of $M=200$ GeV and $M_A=400$ GeV, the cross section is approximately 0.5 pb (or 8 pb) for $\kappa=2$ TeV (or 0.5 TeV), where κ is the new U(1)-symmetry-breaking scale and $M_A=g_A\kappa$ is used. There will thus be a few times 10^3 aspons produced per SSC year.

The decay of the $Q=(U,D)$ will be dependent on the mass of M and we need to consider several cases. It is important that U and D are expected to be almost degenerate in mass, with a splitting $M_U-M_D \approx (\alpha/2\pi)M \approx 10^{-3}M$. Therefore, the decay $U\to DW$ is very suppressed. Thus the decays of quirks have to go through their mixing with usual quarks. The decays are $U\to u^iA$, $U\to d^iW$, $D\to d^iA$, and $D\to u^iW$, where u^i and d^i $(i=1,2,3)$ are the usual quarks. Since W is presumed to be much lighter than A, $U\to d^iW$ and $D\to u^iW$ are dominant.

First, let us consider the U quirk. Since quirks decay democratically into the different light families, we look for the decays $U\to d^iW\to d^il\nu$. Therefore, the signature for $pp\to U\bar{U}$ is two isolated charged leptons and two or more jets with missing energy due to neutrinos. Two isolated charged leptons with different flavor (μ^+e^- or μ^-e^+) help the rejection of background from $b\bar{b}$ production and the Drell-Yan process. However, this signature has the same topology as in the case of $pp\to t\bar{t}$, where t decays into b predominantly. Hence, b-jet identification is important to distinguish between these two processes.

Next, let us consider the D quirk. If $M > m_t+m_W$, the decay $D\to tW\to bWW$ will provide a background-free signal of D. The signature for $pp\to D\bar{D}$ is four isolated charged leptons and two or more jets with missing energy. The corresponding cross section is given by

$$\begin{aligned}\sigma(pp &\to 4l+2\text{ jets}+\text{missing energy})\\ &=\sigma(pp\to D\bar{D})[B(D\to tW)]^2[B(t\to bW)]^2\\ &\qquad\times[B(W\to l\nu)]^4\\ &\approx\sigma(pp\to D\bar{D})(\tfrac{1}{3})^2(1)^2(\tfrac{3}{9})^4.\end{aligned}$$

To obtain 30 events per SSC year requires $\sigma(pp\to D\bar{D})$ to be greater than 1 pb. From Fig. 1, this occurs if $M < 300$ GeV and includes therefore the whole allowed range of M [7]. If $M < m_t+m_W$, we have to look for the decay $D\to u^iW$. This has the same signature as in the case of U quirks.

If M is sufficiently small, $M < 150$ GeV, then $D\bar{D}$ and $U\bar{U}$ could be produced at Fermilab in the top search [6]. The decay topology of U and D decays are then identical to the sought-for t decay, but the Q decays could be distinguished by (i) the degeneracy of U and D, eventually distinguishable by study of the resultant jets, and by (ii) the fact that U and D decay directly into light quarks like d and u, respectively, more frequently than does the t quark.

We turn now to the decay and signature of the aspon. The total decay width Γ_A depends strongly on M_A. If $M_A > 2M$ the dominant decay is $A\to Q\bar{Q}$ with a decay

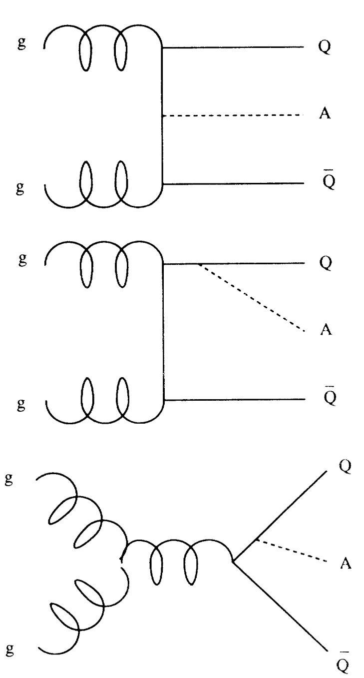

FIG. 2. Feynman diagrams for $pp\to Q\bar{Q}A$ used in the calculation of Fig. 3; further diagrams obtained by crossing must be added.

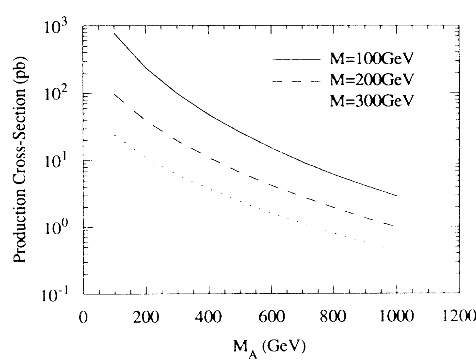

FIG. 3. Cross section for $pp\to Q\bar{Q}A$ from gluon fusion at $\sqrt{s}=40$ TeV and $m=100,200,300$ GeV. $\alpha_A=g_A/4\pi=0.1$ is used.

rate given by

$$\Gamma_A = \alpha_A M_A (1+2M^2/M_A^2)(1-4M^2/M_A^2)^2 ,$$

giving $c\Gamma_A$ of order 1 GeV. If $2M > M_A > M$, the principal decay is $A \to Q\bar{q}_i$, which is suppressed relative to $A \to Q\bar{Q}$ by $|x_i|^2$, where x_i, defined in Ref. [2], is the mixing of q^i and Q; this implies Γ_A of order 1 MeV. In the case $M_A < M$, the decay $A \to q_j \bar{q}_i$ will be of order 1 keV; in this situation the decay of A to three gluons by a box graph is competitive.

Let us consider a typical case $M_A = 400$ GeV and $M = 200$ GeV so that the cross section at SSC for $pp \to Q\bar{Q}A$ is 0.5–8 pb as discussed above, depending on the scale of $U(1)_{\text{new}}$ breaking. This is of similar magnitude to the well-known Higgs-boson production $pp \to t\bar{t}H$ [8], so experiments seeking the Higgs particle can at the same time look for the aspon. Two distinctions between the two processes are that (i) the aspon total width Γ_A is much smaller than the Higgs-boson total width Γ_H, since $\Gamma_A/\Gamma(H \to t\bar{t}) \cong M_A^3 m_W^2/\kappa^2 m_H^3 \le 0.01$, and (ii) unlike $H \to 2\gamma$ [9], $A \to 2\gamma$ is forbidden by Yang's theorem. For a light aspon with $M_A < M$, the decay $A \to q\bar{q}$, where q are in the light families, can give $A \to 2$ jets, which together with the extremely narrow Γ_A will give a distinctive signature.

A promising approach to finding new flavors is through searching for heavy quark bound states. This seems difficult for the $t\bar{t}$ system since the single t quark decay to bW is expected to yield a top lifetime shorter than the formation time of toponia if the t quark is heavier than 100 GeV [10]. However, the single quirk lifetime is expected to be much longer because of the small mixing with three light generations. This implies the existence of bound quirkonium states.

We now discuss the production and decay of quirkonium. At the SSC, the gluon-gluon luminosity is roughly 50 times the $q\bar{q}$ luminosity so gluon fusion is the dominant source of quirkonium production as it is for open flavor production. The production cross section is dominated by η ($J^{PC}=0^{-+}$), ψ (1^{--}), χ_0 (0^{++}), and χ_2 (2^{++}), so we shall restrict our discussion to these modes. The η and $\chi_{0,2}$ states couple to gg and are produced in the gluon fusion reaction. Because ψ states couple only to ggg, the dominant production channel is $gg \to \psi g$.

For quirks in the range allowed by the analysis of Refs. [2] and [5], we have quirkonium mass bounds $200 \le m(Q\bar{Q}) \le 600$ GeV. These bounds correspond to approximate [1] production cross sections for η, ψ, and $\chi_{0,2}$ of $\approx (2\times 10^2)-10^{-3}$ pb, $\approx 1-10^{-2}$ pb, and $\approx (2\times 10^{-1})-10^{-3}$ pb, respectively, or with the design luminosity at the SSC of 10^4 pb^{-1}/yr we expect $(2\times 10^6)-10^4$, 10^4-10^2, and $(2\times 10^3)-10$ events per year, respectively.

Each quirkonium state has approximately $N_s = 2[M_Q/(1.5\text{ GeV})]^{1/2}$ "stable" radial excitations [11]: thus $8 \le N_s \le 28$ here. This forest of excitations leads to an enhancement of the η production cross section by a flat factor of 2 and of χ_0 by a factor of 6 to 4 over the allowed mass range.

For simplicity we assume $M(Q\bar{Q}) < M_A, M_\chi$, where M_χ is the mass of the singlet Higgs scalar. Since the quirks do not couple directly to the standard model Higgs doublet and $M(U)$ is approximately degenerate with $M(D)$, the dominant quirkonium decay modes are through

(a) $\vartheta_{Q\bar{Q}} \to \gamma^* \to p_1 p_2$,

(b) $\vartheta_{Q\bar{Q}} \to Z^* \to p_1 p_2$,

(c) $\vartheta_{Q\bar{Q}} \to t$-channel fermion $\to p_1 p_2$.

We list the three classes of dynamically allowed two-body decay modes of quirkonium in the following table, where $\vartheta_{Q\bar{Q}}$ stands for any quirkonium:

			$\vartheta_{Q\bar{Q}}$	
p_1p_2	$\eta(^1S)$	$\psi(^3S)$	$\chi_0(^3P)$	$\chi_2(^3P)$
$f_1\bar{f}_2$	...	a,b	...	...
W^+W^-	c	a,b,c	c	c
$gg,\gamma\gamma,ZZ$	c	...	c	c
$Z\gamma$	c	c	c	c

The production cross section of open flavor quirks at the SSC is larger than that of quirkonium and here would be the discovery mode; quirkonium might subsequently be discovered at the SSC by its $\eta \to W^+W^-$ or $\psi \to l^+l^-$ decay modes. At the proposed Next Linear Collider (NLC), a detailed study of quirkonia would be feasible including the spectrum, decay widths, and branching ratios. The lifetimes and branching ratios are different from the quarkonia counterparts because the weak interactions are very different [2]; the interquirk potential, on the other hand, is similar (aside from the overall mass scale) to the interquark potential [12].

In conclusion, we have computed the production cross sections of $pp \to Q\bar{Q}$ and $Q\bar{Q}A$ in hadron colliders. The signature of $pp \to Q\bar{Q}$ is two isolated opposite-charged leptons and two or more jets with missing energy due to neutrinos. This signature is topologically identical to $pp \to t\bar{t}$, but Q decays democratically to the light quarks while t decays dominantly to b. Therefore, b-jet identification is important to distinguish between these two processes. For heavy quirks, $M > m_t + m_W$, the background-free signature for $pp \to Q\bar{Q}$ is four isolated charged leptons and two or more jets with missing energy. Requiring 30 events per SSC year, we find that the D quirk will be discovered if its mass is less than 300 GeV. Bare Q decay is highly suppressed relative to the rapid $t \to bW$ and hence $Q\bar{Q}$ states have ample time to form quirkonium. A detailed study of quirkonia would be feasible at the NLC. The aspon, fairly independently of its mass, is expected to decay into two jets. The signature will be typically identical to $pp \to t\bar{t}H$. Therefore, experiments seeking $pp \to t\bar{t}H$ can at the same time look for the aspon.

Finally, it is important to emphasize the generic qualities of the aspon model in that one always requires a vector aspon and quirk. There is very little room (much less than anticipated in Ref. [1]) to extend or embed the aspon model in any satisfying way. The aspon model makes testable phenomenological predictions and we must anxiously await the data.

We thank T. Rizzo for providing the parton structure functions. This work was supported in part by the U.S. Department of Energy under Grant No. DE-FG05-85ER-40219 and Grant No. DE-FG05-85ER-40226.

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