served in metal films remains. The evidence (particularly the spatial cross-correlation experiment[1]) presented by Voss and Clarke that this noise results from "temperature" fluctuations remains convincing. The arguments of Horn and co-workers[5-7] that the noise magnitude does not have the right temperature dependence to be the predicted spontaneous enthalpy fluctuations are also convincing and agree with our conclusion that the spontaneous noise should be scarcely measurable in ordinary metal films. The conclusion of Voss and Clarke that the spectrum of the spontaneous fluctuations is skewed by unknown forces[4] appears to be unjustified, as is the conclusion of Horn and co-workers that the predominant source of the $1/f$ noise is not temperature fluctuations at all.[6,7] An intermediate conclusion which seems consistent with all the data would be that the $1/f$ noise comes from "temperature" fluctuations of unknown origin in excess of those predicted by thermodynamics.

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$$\tau_H = (\Delta R/R)c_p WT(P\gamma)^{-1} = S(0)/4\int_0^\infty S(f)\,df$$

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Problem of $D^0 \to K^-\pi^+\pi^0$ in Nonleptonic Charm Decay

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If the branching ratio for $D^0 \to K^-\pi^+\pi^0$, which currently stands at $(12 \pm 6)\%$, remains much larger than 3% in future data, then either the $\Delta T = 1$ rule breaks down, or the final-state $\bar{K}\pi\pi$ contains much structure. The possibility that this structure is due to $\bar{K}^*$ and ρ resonances is briefly examined.

Recent data[1-3] on $D \to \bar{K}\pi\pi$ are examined in the light of the $\Delta T = 1$ rule for nonleptonic charm decay.[4] It is argued that if the branching ratio for $D^0 \to K^-\pi^+\pi^0$, which currently stands at $(12\pm 6)\%$,[2] remains significantly larger than 3%, then the $\bar{K}\pi\pi$ final state must be a highly structured object, and that this structure could well be due to resonances such as $\bar{K}^*$ and ρ. Implications for the ratio of the total D^+ and D^0 widths and for $D \to \bar{K}\pi$ are also considered.

Although the observed sample of $D \to \bar{K}\pi\pi$ is small, it seems to indicate that the corresponding decay matrix element is uniform over the Dalitz plot.[3] If this is indeed the case, then the $\Delta T = 1$ rule predicts a simple relationship[5] between the rates for $D^+ \to K^-\pi^+\pi^+$ and $D^0 \to K^-\pi^+\pi^0$:

$$\Gamma(K^-\pi^+\pi^+) = 4\Gamma(K^-\pi^+\pi^0), \tag{1}$$

where $\Gamma(K^a\pi^b\pi^c)$ denotes the rate for a specific decay mode of the appropriate D meson. According to recent measurements, the branching ratios for these decay modes are (3.9 ± 1.0)% and (12 ± 6)% respectively (see Table I). Consequently the ratio of the total widths of D mesons can be deduced from Eq. (1):

$$R \equiv \Gamma(D^+)/\Gamma(D^0) \approx 12 \pm 6, \tag{2}$$

where $\Gamma(D^x)$ denotes the total width of D^x, and 3.9 ± 1.0 has been set equal to 4 for the purpose of this discussion.

Now it has been shown by Peshkin and Rosner,[6] and by Pais and Treiman,[7] that the $\Delta T = 1$ rule imposes strict bounds upon the ratio of widths:

$$0 \leq R \leq 3. \tag{3}$$

Although the branching ratio for $D^0 \to K^-\pi^+\pi^0$ is not accurately measured at present, a comparison of Eqs. (2) and (3) indicates that there may be an apparent conflict with the $\Delta T = 1$ rule. If the experimental number (12 ± 6)% is found to be less than 3% on subsequent experiments, then there need be no problem; but if (12 ± 6)% turns out to be greater than 3%, then the value of R as determined from Eq. (1) will definitely contradict the upper bound of Eq. (3).

One way of resolving such a conflict would be to give up the $\Delta T = 1$ rule. This rule, however, is a basic feature of the standard charm model,[4] and so it is desirable to maintain its validity as long as possible. Therefore it may become necessary to modify the other assumption upon which Eq. (1) is based, namely the uniformity of the $D \to \bar{K}\pi\pi$ matrix element over the Dalitz plot.

As remarked above, there is no positive evidence for nonuniformity in the present data[3]; but the number of events is sufficiently small, being less than 150, that some lack of uniformity could begin to emerge as more events are collected. Indeed, if the $\Delta T = 1$ rule is to be preserved, then the experimental alternatives at this time are either that the branching ratio for $D^0 \to K^-\pi^+\pi^0$ be less than 3%, or that the decay matrix element have a nonuniform structure over the Dalitz plot. Ultimately it may transpire that both alternatives are found to be valid.

In the event that it becomes necessary to introduce structure into the Dalitz plot for $D \to \bar{K}\pi\pi$, one obvious way of doing so is to suppose that the final state contain substantial components of $\bar{K}^*\pi$, or $\bar{K}\rho$, or both.[4,8] The final state is then fed by three channels: a direct decay with a uniform matrix element; a decay into $\bar{K}^*(890)\pi$ followed by $\bar{K}^* \to \bar{K}\pi$; and a decay into $\bar{K}\rho$ followed by $\rho \to \pi\pi$. In general there are six independent amplitudes, two from each channel, with which to describe the five charge states of $D \to \bar{K}\pi\pi$, and so there is more than enough freedom to fit any set of branching ratios without contradicting the $\Delta T = 1$ rule.

The number of amplitudes can be reduced by assuming that the Cabibbo-allowed nonleptonic Hamiltonian is dominated by the sextet representation of SU(3).[4] In this case the direct channel is still described by two amplitudes, but the resonant channels are now described by three amplitudes instead of four, as shown in Table II. There is still considerable freedom in fitting the branching ratios, as can be exemplified by some special cases.

As one example, consider the special case in which the amplitudes for $\bar{K}^*\pi$ vanish and the only resonant channel is $\bar{K}\rho$; this corresponds to setting

$$\boldsymbol{T} = S + \boldsymbol{A} = 0 \tag{4}$$

in Table II. If interference effects between $\bar{K}\rho$ and the direct $\bar{K}\pi\pi$ channel are negligible, and if (see Table I)

$$\begin{aligned} &\Gamma(K^-\pi^+\pi^+) \approx (4\%) \times \Gamma(D^+), \\ &\Gamma(\bar{K}^0\pi^+\pi^-) \approx (4\%) \times \Gamma(D^0), \\ &\Gamma(K^-\pi^+\pi^0) \approx (x\%) \times \Gamma(D^0), \end{aligned} \tag{5}$$

then the following relations can be derived from Eq. (4) and Tables I and II:

$$\begin{aligned} &\Gamma(\bar{K}^0\rho^+) = 0, \\ &2\Gamma(\bar{K}^0\rho^0) = \Gamma(K^-\rho^+) = [(x - R)\%] \times \Gamma(D^0), \\ &8 - (x - R) = \tfrac{4}{9} R\,|1 + (2C_0/C_2)|^2, \end{aligned} \tag{6}$$

where $\Gamma(K^a\rho^b)$ denotes the rate for the specific decay mode of the appropriate D meson, and R is the ratio of total widths [see Eq. (2)]. Since the right-hand sides of Eqs. (6) cannot be negative, it follows that

$$8 \geq (x - R) \geq 0. \tag{7}$$

TABLE I. Branching ratios for D-meson decay.

Decay mode	Branching ratio (%)	Reference
$K^-\pi^+$	2.2 ± 0.6	1
$\bar{K}^0\pi^+\pi^-$	4 ± 1.3	1
$K^-\pi^+\pi^0$	12 ± 6	2
$\bar{K}^0\pi^+$	1.5 ± 0.6	1
$K^-\pi^+\pi^+$	3.9 ± 1.0	1

TABLE II. Amplitudes for $D \to \bar{K}\pi\pi$. The direct decay amplitudes C_0 and C_2 correspond to 2π final states with $T=0$ and 2, respectively. The resonant amplitudes are taken from Einhorn and Quigg (Ref. 4).

Mode	Direct	$\bar{K}^{*0}$	K^{*-}	ρ
$K^-\pi^+\pi^+$	$3C_2$	$-\sqrt{6}T$	0	0
$\bar{K}^0\pi^+\pi^0$	$3C_2/2$	$\sqrt{3}T$	0	$-3T$
$K^-\pi^+\pi^0$	$3C_2/2$	$-\frac{1}{3}\sqrt{3}(2T-S-A)$	$-\frac{1}{3}\sqrt{3}(T+S+A)$	$-T+S-A$
$\bar{K}^0\pi^+\pi^-$	$\sqrt{2}C_0+C_2/\sqrt{2}$	0	$\frac{1}{3}\sqrt{6}(2T+S+A)$	$-\frac{1}{2}\sqrt{2}(2T+S-A)$
$\bar{K}^0\pi^0\pi^0$	$-C_0+C_2$	$\frac{1}{6}\sqrt{6}(2T-S-A)$	0	0

Thus, depending on the value of R [see Eq. (3)], the branching ratio for $D^0 \to K^-\pi^+\pi^0$ could range anywhere from zero to 11%.

Another possibility is to assume that $\bar{K}^*\pi$ is the only resonant channel, and that the amplitudes for $\bar{K}\rho$ in Table II vanish: i.e.,

$$T=S-A=0. \tag{8}$$

In this case, it follows that

$$\begin{aligned} 2\Gamma(\bar{K}^{*0}\pi^0) &= \Gamma(K^{*-}\pi^+) = [\tfrac{3}{2}(x-R)\%]\times\Gamma(D^0), \\ 4-(x-R) &= \tfrac{2}{9}R\,|1+(2C_0/C_2)|^2, \end{aligned} \tag{9}$$

as long as interference between different channels can be ignored. Both sides in Eq. (9) must be positive, and so

$$4 \geq (x-R) \geq 0. \tag{10}$$

Thus, if $\bar{K}^*\pi$ is the only resonant channel, then x lies somewhere in the range 0–7%.

Equations (4) and (8) have two consequences in common: first that the rate for $D^0 \to \bar{K}^0\pi^+\pi^0$ is given by

$$\Gamma(\bar{K}^0\pi^+\pi^0) = \tfrac{1}{4}\Gamma(K^-\pi^+\pi^+) \approx (1\%)\times\Gamma(D^+), \tag{11}$$

and second that there are no resonances in the D^+ decay. Thus a relatively low rate for $D^+ \to \bar{K}^0\pi^+\pi^0$ as in Eq. (11) together with the occurrence of either a $\bar{K}^*$ or ρ in D^0 decay serves as tests of the possiblity that $D \to \bar{K}\pi\pi$ arises from a direct channel plus a single resonant channel.

It may happen that the final state is dominated by resonances and that the direct channel is suppressed:

$$C_0 \approx C_2 \approx 0. \tag{12}$$

If interference between different resonances is neglected and phase-space factors for $\bar{K}^*\pi$ and $\bar{K}\rho$ are taken to be roughly equal, then the rate for $D^+ \to \bar{K}^0\pi^+\pi^0$ is large:

$$\Gamma(\bar{K}^0\pi^+\pi^0) \approx 2\Gamma(K^-\pi^+\pi^+) \approx (8\%)\times\Gamma(D^+). \tag{13}$$

In addition, the ratio of rates for the $K^-\pi^+\pi^0$ and $K^-\pi^+\pi^+$ decays [see Table II and Eqs. (5) and (12)] gives rise to a lower bound for x,

$$x \geq R, \tag{14}$$

but no upper bound. Other branching ratios will depend upon R, the ratio of total widths.

The parameter R plays an important role in another aspect of nonleptonic charm decay, namely the question of sextet dominance in $D \to \bar{K}\pi$. Sextet dominance requires that the amplitude for $D^+ \to \bar{K}^0\pi^+$ be much smaller than the amplitude for $D^0 \to K^-\pi^+$.[4] Recent measurements indicate that the branching ratios of these two decay modes are comparable with one another, thus the amplitude for the D^+ decay mode will be smaller than that for the D^0 mode only if the total D^+ width is smaller than the total D^0 width. It will therefore be very interesting to determine R from the ratio of semileptonic decay rates.[9]

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Note added.—As this work was going to press, the author learned that S. Kaptanoglu (SLAC-PUB 2072) has reached similar conclusions using different arguments.

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Is the Υ a Bound State of Exotic Quarks?

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To explain the Υ-Υ′ mass splitting of the recently discovered Υ resonances, we interpret them as bound states of a quark and an antiquark which are not members of the fundamental (3) representation of color SU(3). A quark in the 8 or 6 representation is consistent with the data. We discuss some of its properties and the phenomenological consequences that follow from our suggestion.

It is widely accepted that the newly discovered upsilon resonance[1] Υ(9.4) is a bound state ($b\bar{b}$) of a fifth quark b and its antiquark. However, the experimentally observed splitting[2] of the Υ and Υ′ is too big to be explained by the same (essentially) linear potential that has been applied so successfully to the phenomenology of the ψ spectroscopy[3] (at least in the nonrelativistic approximation). To be specific, the potential used by Eichten *et al.*[4] gives a splitting of 0.4 GeV instead of the observed $M_{\Upsilon'} - M_{\Upsilon} \simeq 0.6$ GeV.

To explain the observed splitting, other nonrelativistic potential models have been suggested.[5,6] For example, a logarithmic potential[5] seems to be consistent with both the ψ and Υ spectroscopies in the nonrelativistic limit. However, such potential models provide little connection to other aspects of hadronic phenomena.

On the other hand, recent theoretical investigations on quark confinement, linear Regge trajectories for light hadrons, ψ spectroscopy, and other hadronic properties strongly indicate a stringlike (or vortexlike) picture for hadrons. A linear potential arises nonrelativistically in any stringlike model. Hence, before abandoning such an intuitively simple picture, it is important to find an explanation for the large Υ-Υ′ mass splitting within a simple string model.

One obvious possibility is that Υ and Υ′ are bound states of two different types of quarks. In particular, in either a simple linear potential[7] or the modified one of Ref. 4, the first radially excited state of the ground state Υ(10.0) occurs at roughly 10.4 GeV, precisely where some structure,[2] namely Υ″(10.4), has been observed. The real Υ′ is then presumably buried between the Υ(9.4) and Υ(10.0).

Another possible explanation within the string picture, suggested by Giles and Tye,[8] is that the Υ (and Υ′) is a bound state of an exotic quark-antiquark pair, i.e., quarks which are not in the fundamental representation 3 of color SU(3).[9,10] In this Letter we would like to discuss the likelihood of such an exotic quark. We elaborate on the dynamical reasoning behind this possibility and discuss some of its fascinating phenomenological consequences.

Since hadron physics is still poorly understood, our dynamical reasoning must necessarily be model dependent. We shall assume exact color confinement; for the sake of clarity and definiteness we shall carry out our discussion within the framework of the quark-confining string (QCS) model,[11] which is assumed here to be a phenomenological model[12,13] of quantum chromodynamics (QCD). Our attitude and philosophy are very similar to those of Eichten *et al.*[4] Hence it should be obvious that the validity of our argument is more

