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# Do the $P_c^+$ Pentaquarks Have Strange Siblings?

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The recent LHCb discovery of states $P_c^+(4380)$, $P_c^+(4450)$, believed to be $c\bar{c}uud$ pentaquark resonances, begs the question of whether equivalent states with $c\bar{c} \to s\bar{s}$ exist, and how they might be produced. The precise analogue to the $P_c^+$ discovery channel $\Lambda_b \to J/\psi\, K^-p$, namely, $\Lambda_c \to \phi\pi^0 p$, is feasible for this study and indeed is less Cabibbo-suppressed, although its limited phase space suggests that evidence of a $s\bar{s}uud$ resonance $P_s^+$ would be confined to the kinematic endpoint region.

## I. INTRODUCTION

The discovery of multiple exotic charmoniumlike (and bottomoniumlike) states in the past dozen years, starting with Belle's observation [1] of the $X(3872)$, has been nothing short of stunning. Previously, the heavy quarkonium $Q\bar{Q}$ systems were the best understood sectors of hadronic physics, with spectra that could be completely predicted from two-body Schrödinger equations, modeled using nonrelativistic potentials developed decades ago and continuously refined [2–4]. The observation of more than 20 such so-called $X, Y, Z$ bosonic states [5] (several of which have been confirmed by multiple experiments) has upended this previous simple picture. The latest finding in this regard is the LHCb observation [6] of two exotic baryonic charmoniumlike states, $P_c^+(4380)$ and $P_c^+(4450)$, at high statistical significance in the $J/\psi\, p$ spectrum of $\Lambda_b \to J/\psi\, K^-p$. These states, like the $Z(4475)$ [also called $Z(4430)$] before them [7], are shown to have rapid phase variation in their production amplitudes consistent with true resonant behavior. Evidence continues to mount that at least some of $X$, $Y$, and $Z$ are genuine $c\bar{c}q_1\bar{q}_2$ tetraquark states (not just kinematical effects), while the $P_c^+$ are $c\bar{c}uud$ pentaquark states.

Of course, the possibility of QCD exotics has been noted in the earliest days of the quark model [8, 9], even before the advent of color dynamics. Nevertheless, despite decades of scrutiny, no unambiguous experimental signal indicating the existence of an exotic hadron outside of the $q\bar{q}$-meson, $qqq$-baryon paradigm has ever been identified in the light-quark $(u, d, s)$ sector, or even in the sector with a single heavy quark. Why should it be that exotics are first becoming visible in doubly heavy-quark systems? One explanation lies in the embarrassment of hadronic riches in the $< 2.5$ GeV range: Many of the purported light-quark exotics have the same quantum numbers as conventional quark-model states and can hide amongst them, or indeed, mix with them quantum mechanically. Even the extremely well-established $J^{PC} = 1^{++}$ $X(3872)$ might mix at some level with the yet-unseen conventional $c\bar{c}$ state $\chi_{c1}(2P)$ (as suggested numerous times, most recently in Ref. [10]).

Previous work by the present author [11–15] to explain this curious fact has argued that the key feature in forming an identifiable exotic state is the presence of two components in the hadron, each of which contains a heavy quark and is consequently fairly compact (a few tenths of a fm), but that are separated from each other, in the sense of having a small wave function overlap, by a somewhat larger distance. The specific proposal in those works is the presence of compact colored diquark (or triquark) components separately bound together through the attractive $\mathbf{3} \otimes \mathbf{3} \supset \bar{\mathbf{3}}$ color interaction and collectively bound together by confinement. However, the same situation arises in the molecular picture, in which the compact components are color-singlet heavy quark-containing meson and baryon pairs. In the former case, the residual color interaction between the components is full-strength QCD, and in the latter it is the much weaker residual color van der Waals QCD force.

Supposing that exotics have only recently become observable because they require well-separated components, each containing a heavy quark, one may reconsider analogous systems in which the $c\bar{c}$ or $b\bar{b}$ quark pairs are replaced by $s\bar{s}$. This proposal is not at all guaranteed success, since the $s$ is not truly heavy (*i.e.*, the current quark mass $m_s$ is smaller than $\Lambda_{\rm QCD}$), and it may well turn out that exotics of the type thus far seen absolutely require the presence of two heavy quarks. A major thrust of this work is to provide one test of this point of view: Is the divide between the $c$ and $s$ quarks so great that *no* exotic behavior survives in the $s\bar{s}$ sector? In order to do so, one may seek out effects perhaps not as prominent as in the heavy-quark systems since $s$ is not truly heavy, but anomalous nonetheless. This proposal was first advocated and applied to the case of $\phi$-$N$ photoproduction in Ref. [15], where it was used to explain the appearance in CLAS (JLab) data [16, 17] of peculiar enhancements of the $\gamma p \to \phi p$ cross section in the $\phi$ forward and backward directions. It was argued that treating the process as a $2 \to 2$ scattering resulting from the formation of a color-antitriplet $(su)$ diquark and a color-triplet $[\bar{s}(ud)]$ antitriquark, which subsequently hadronize to $\phi$ $[= (\bar{s}s)]$ and $p$ $[= u(ud)]$ through the large-separation wave function

tails of the hadrons stretching between the two compact colored components, provides a natural mechanism to explain forward and backward cross section enhancements. In this case, the anisotropic nature of the enhancements disfavors a resonant origin [17], meaning that the $(su)$-$[\bar{s}(ud)]$ complex is only a "would-be" pentaquark [15].

The original $\phi$ photoproduction study was suggested by the simple substitution of $c\bar{c} \to s\bar{s}$ into the $\gamma N \to P_c \to J/\psi N^{(*)}$ photoproduction proposals of Refs. [18–20]. In both $\phi$ and $J/\psi$ cases, the $Q\bar{Q}$ pair arises through the dissociation of the incoming photon. No particular $P_c$ compositeness substructure (diquark [21–25] or other colored substructure [26], molecular [27–32],[1] hadrocharmonium [19], or soliton [40]) is presupposed; but if the states turn out to be kinematical effects [41–43] due to the particular placement of hadronic levels with respect to the original $\Lambda_b \to J/\psi K^- p$ production process, then one would not necessarily expect interesting structure to arise in $\gamma N \to J/\psi N^{(*)}$. A discussion of the relative merits of various interpretations for $P_c^+$ states appears in Ref. [44].

Investigations have also begun into decays related to $\Lambda_b \to J/\psi K^- p$, such as via $\Xi_b$ and $\Omega_b$ [45–47], in a search for the flavor SU(3) partners of the $P_c^+$ states. In each case, the underlying weak decay is $b \to cW^{*-} \to c\bar{c}s$ or $b \to cW^{*-} \to c\bar{c}d$, meaning that the relevant Cabibbo-Kobayashi-Maskawa (CKM) matrix element combination is $V_{cb}V_{cs}^*$ or the even smaller combination $V_{cb}V_{cd}^*$. The possibility that the weak decay is actually $b \to u\bar{u}s$ is explored in Ref. [48].

In this short paper we propose to create the precise hidden-strangeness analogues $P_s^+ = s\bar{s}uud$ to the $P_c^+$ states through the decay $\Lambda_c \to P_s^+ \pi^0 \to \phi \pi^0 p$. The replacement of $b \to c$, specifically the substitution of the weak decay $c \to sW^{*+} \to s\bar{s}u$, is all that is needed to produce this channel from one already known. We illustrate in Fig. 1 the flow of quark flavors in the process in terms of the diquark-triquark picture used for $P_c^+$ in Ref. [14]; however, as argued above, the diquark picture is not necessarily the only one that produces viable double-heavy pentaquark states, and the figure is intended merely as an illustration of one particular viable physical process. Again, we emphasize that $s\bar{s}$-containing exotics, and the $P_s^+$ states in particular, are not guaranteed to exist; the proposal here is that the underlying mechanism creating the $P_c^+$ states also holds for $P_s^+$, and the sole dynamical input is the assumption that the gluodynamics leading to the energy differences involved is flavor independent.

A discussion of the merits and drawbacks of this and related modes is presented in Sec. II, followed by a brief summary in Sec. III.
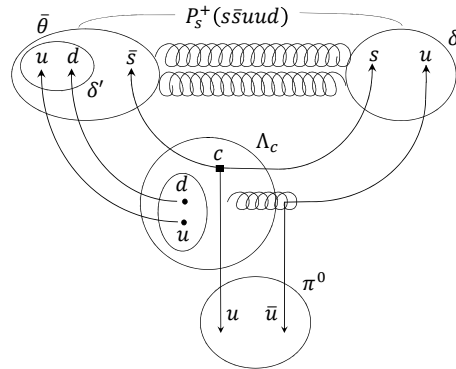


FIG. 1: Illustration of one particular mechanism [diquark ($\delta$)-antitriquark ($\bar{\theta}$) formation] for the production of a pentaquark $P_s^+ = s\bar{s}uud$ state in the decay $\Lambda_c \to P_s^+ \pi^0$. The black square indicates the $c$-quark weak decay. The spatially extended $\delta\bar{\theta}$ state is held together by long-range color forces (indicated by gluon lines) via a color flux tube.

## II. DECAY MODES FOR PRODUCING HIDDEN-STRANGENESS PENTAQUARKS

We begin with the analogy between the process $\Lambda_b \to P_c^+ K^- \to J/\psi K^- p$ and $\Lambda_c \to P_s^+ \pi^0 \to \phi \pi^0 p$, the former illustrated in Ref. [14] and the latter here in Fig. 1. Just as for the corresponding $P_c^+$ production process in Ref. [14], the presence of more than one copy of the light $u$ quark in the final state (indeed, three $u$'s for $P_s^+ \pi^0$ compared to only two for $P_c^+ K^-$) offers the possibility of several production diagrams, regardless of the physical process. In the diquark picture, it is natural to expect the $u$ in the diquark $\delta'$ contained in the initial $\Lambda_{b,c}$ to maintain its identity throughout the process as a spectator. However, the $u$ appearing in the diquark $\delta$ can still emerge either from the weak decay $c \to s\bar{s}u$ or from the $u\bar{u}$ production. Analogous statements apply to other physical decay pictures.

Replacing $u\bar{u} \to d\bar{d}$ produces the isospin-partner process $\Lambda_c \to P_s^0(s\bar{s}udd)\pi^+ \to \phi\pi^+ n$. The rate for this process should be comparable to that for $P_s^+$, but not exactly equal (even in the limit of perfect isospin), since the $c \to s\bar{s}u$ weak decay violates isospin, failing to produce a $d$ quark able to interfere with the ones from other sources.[2] In any case, neutron reconstruction can be experimentally rather challenging, so the $P_s^+$ channel seems

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[1] For predictions of hidden-charm baryons prior to the $P_c^+$ observation, see Refs. [33–39].

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[2] The $d$ quark reemerges if one considers the weak decay $c \to d\bar{s}u$, but such processes are not only CKM suppressed (by a factor $|V_{cd}^*/V_{cs}^*|^2 \simeq \frac{1}{20}$), but also may lack an $s$ quark in the final state.

to be the most promising one for near-term investigations.

To be precise, if the mixing associated with the weak-decay $u$ quark is ignored, the processes $\Lambda_b \to P_c^+ K^- \to J/\psi\, K^- p$ and $\Lambda_c \to P_s^+ \pi^0 \to \phi\, \pi^0 p$ are entirely comparable if one substitutes $V_{cb}V_{cs}^* \to V_{cs}^* V_{us}$ and notes that $K^-$ and $\pi^0$ are SU(3)-partnered pseudo-Nambu–Goldstone bosons. Only the $\frac{1}{\sqrt{2}}$ in the $\pi^0$ flavor wave function is different.

Noting that $|V_{cb}V_{cs}^*|^2/|V_{cs}^* V_{us}|^2 \simeq \frac{1}{30}$, one sees that the corresponding charmed decays are substantially less CKM suppressed than their bottom counterparts. This simple fact explains why many of the lighter $b$-hadrons have lifetimes comparable to their charmed counterparts despite a much greater available phase space. Conversely, the branching fractions for the corresponding charmed processes such as $\Lambda_c \to P_s^+ \pi^0 \to \phi\, \pi^0 p$, or for that matter, nonresonant $\Lambda_c \to \phi\, \pi^0 p$, are enhanced compared to their $b$ counterparts. A glance at the known branching fractions for $\Lambda_c$ decays [49] shows known three-body decays to occur in the several times $10^{-3}$ range or even larger, indicating that, if nothing else, the discovery of the mode $\Lambda_c \to \phi\, \pi^0 p$ should be straightforward.

The resonance $P_s^+$ content of the process is also expected to be nonnegligible. In the $P_c^+$ case, LHCb measured [6] the branching fraction ratios

$$\frac{\text{B.R.}(\Lambda_b \to P_c^+ K^-)}{\text{B.R.}(\Lambda_b \to J/\psi K^- p)} = \begin{cases} (8.4 \pm 0.7 \pm 4.2)\% \\ (4.1 \pm 0.5 \pm 1.1)\% \end{cases}, \quad (1)$$

for $P_c^+(4380)$ and $P_c^+(4450)$, respectively, quite significant considering the large available $\Lambda_b$ phase space, suggesting that $P_s^+$ production in $\Lambda_c \to \phi\, \pi^0 p$ will not be uncommon.

Arguably, the most interesting difference between $\Lambda_b \to P_c^+ K^- \to J/\psi\, K^- p$ and $\Lambda_c \to P_s^+ \pi^0 \to \phi\, \pi^0 p$ is simple phase space. Since $m_b - m_c$ is so much greater than $m_c - m_s$, a much greater phase space is available in the former process. To be specific,

$$\begin{aligned} m_{\Lambda_b} - m_{J/\psi} - m_p - m_{K^-} &= 1090.64 \pm 0.23 \text{ MeV}, \\ m_{\Lambda_c} - m_\phi - m_p - m_{\pi^0} &= 193.75 \pm 0.14 \text{ MeV}. \end{aligned} \quad (2)$$

Note first that the phase space for the $\Lambda_c$ decay is so small that no unflavored meson in the final state heavier than $\pi$, and no unflavored baryon heavier than a nucleon, is possible. In the case of final-state decays to states in which the heavier quarks emerge in separate hadrons:

$$\begin{aligned} m_{\Lambda_b} - m_{\overline{D}^{*0}} - m_{\Lambda_c} - m_{K^-} &= 832.40 \pm 0.28 \text{ MeV}, \\ m_{\Lambda_c} - m_{K^{*+}} - m_\Lambda - m_{\pi^0} &= 144.14 \pm 0.29 \text{ MeV}, \end{aligned} \quad (3)$$

which assumes that the mesons appearing in the $P_{c,s}^+$ resonance decays preferentially have $J^P = 1^-$ like $J/\psi$ or $\phi$, then the phase space is even smaller. Obviously, if $\bar{D}^0$ or $K^+$ final states can occur in the $P_c^+$ or $P_s^+$ decays,

respectively (which requires higher partial waves since the $P_c^+$ states have $J = \frac{3}{2}$, $\frac{5}{2}$), then the phase space is correspondingly larger. But the message is clear: If $P_s^+$ resonances are formed in the decay of $\Lambda_c$, they do not have much available phase space.

This effect is further magnified if one notes that the observed $P_c^+$ resonances lie well above the $J/\psi\, p$ threshold, by about 415 and 345 MeV for $P_c(4450)$ and $P_c(4380)$, respectively. If similar numbers hold for the distance of the purported $P_s^+$ states from the $\phi\, p$ threshold, which assumes a flavor independence of the mechanism depicted in Fig. 1 (an extremely crude first approximation), then according to Eqs. (2)–(3), the resonance peaks will not be visible in $\Lambda_c$ decays. Indeed, the $P_c(4450)$ is sufficiently narrow ($\Gamma = 39 \pm 20$ MeV) that a corresponding $P_s^+$ resonance (peak at 2372 MeV) would not be visible at all. However, the large width of $P_c(4380)$ ($\Gamma = 205 \pm 88$ MeV) suggests that an exact $P_s^+$ analogue (peak at 2303 MeV) would begin to appear in the endpoint region. The topology of the event in the center-of-momentum (c.m.) frame of the $\Lambda_c$ in such a case would be remarkable: Since this c.m. frame is the rest frame of not only the $\Lambda_c$ but almost that of the $P_s^+$ and $\pi^0$ as well, one would observe two photons of nearly equal energy from the $\pi^0$ decay emerge back-to-back, as well as a slowly moving proton recoiling against a nearly collinear $K\bar{K}$ pair from the $\phi$ decay.

The higher-strangeness SU(3) partners to the decay $\Lambda_c \to \phi\, \pi^0 p$ whose parent baryon decays weakly, namely $\Xi_c^+ \to \phi\, \overline{K}^0 p$ and $\Omega_c \to \phi\, \overline{K}^0 \Lambda$, tend to have even less phase space: about 13 and 62 MeV, respectively. However, the decay $\Xi_c^+ \to \phi\, \pi^+ \Lambda$ has almost precisely the same phase space as $\Lambda_c \to \phi\, \pi^0 p$:

$$m_{\Xi_c^+} - m_\phi - m_\Lambda - m_{\pi^+} = 193.22 \pm 0.40 \text{ MeV}. \quad (4)$$

Not only does the $\Xi_c^+$ have a longer lifetime than $\Lambda_c$ ($\tau_{\Xi_c^+} = 0.44$ ps, to be compared with $\tau_{\Lambda_c} = 0.20$ ps or $\tau_{\Lambda_b} = 1.47$ ps), it has a perhaps more easily reconstructed final state: a primary decay $\pi^+$ and a secondary decay $\Lambda \to p\pi^0$. Note that the channels with $\phi\Lambda$ in the final state imply open-strangeness ($s\bar{s}uds$) pentaquarks, which may be expected to lie above $P_s^+ = s\bar{s}uud$ states in mass, and therefore potentially outside of range of the initial hadron phase space. It is also worth noting that the $\Xi_c^+ = cus$ and $\Omega_c = css$ contain diquarks of a somewhat different nature than the $(ud)$ in a $\Lambda$ state, the former carrying nonzero isospin and the latter requiring proper antisymmetrization between the identical $s$ quarks. Indeed, part of the motivation for this work (see Fig. 1) was to treat the $s$ quarks as heavy, which becomes less compelling if $s$ quarks arise elsewhere in the process.

The small available phase space remains the least appealing feature of this proposal. Even so, several points are worth mentioning. First, the $P_s^+$ mass estimates may be unnecessarily high. Since the observed $P_c^+$ states have the large spins $J = \frac{3}{2}$ and $\frac{5}{2}$ (as well as one of them necessarily having negative parity), it is very likely that lighter

yet-unobserved $P_c^+$ states exist, which implies lighter $P_s^+$ would be possible as well. Indeed, no obvious physical principle requires the $P_s^+$ states to lie above the $\phi\,p$ threshold as far as the corresponding $P_c^+$ states lie above the $J/\psi\,p$ threshold. Second, the large available phase space for $P_c^+$ decay was actually a substantial nuisance in the $P_c^+$ observation paper [6], since multiple excited $\Lambda$ states had to be included in the analysis. With phase space as small as suggested above, potential higher states provide very little contamination.

One may hope to avoid the phase-space problem, as well as reduce the CKM suppression, by considering not $c \to sW^{*+} \to s\bar{s}u$ but $c \to sW^{*+} \to s\bar{d}u$, which increases the production rate of relevant processes by a factor of $|V_{ud}/V_{us}|^2 \simeq 20$. However, the $P_s^+$ state still requires an $\bar{s}$ quark, which must now appear through pair production. The final state then requires one to accommodate at least three strange quarks, creating insurmountable difficulties with phase space in charmed baryon decays. Alternately, one may attempt to build states in the manner of Fig. 1 without an $\bar{s}$ quark in the antitriquark $\bar{\theta}$, but doing so violates the premise of exotics requiring a heavy quark in each component to be detectable.

## III. CONCLUSIONS

A search for the hidden-strangeness pentaquarks $P_s^+ = s\bar{s}uud$, siblings to the newly observed $P_c^+$ pentaquark candidates, appears to be well within current experimental capabilities. The main difficulty appears to be the limited phase space available in the decays of the likely charmed baryon sources, the best candidate being $\Lambda_c \to P_s^+\pi^0 \to \phi\,\pi^0 p$, which is the exact SU(3)-flavor analogue to the channel $P_c^+$ discovery channel $\Lambda_b \to J/\psi\,K^- p$. The decay $\Xi_c^+ \to \phi\Lambda\pi^+$ appears to be interesting both for the relative ease of its reconstruction and the possibility of finding hints of an open-strangeness $s\bar{s}uds$ pentaquark. At minimum, new unobserved decay modes are within reach, but with a bit of luck, more exciting hints of exotic hadron structure may be uncovered.

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