Crystal ball. This particle detector consists of 672 large sodium iodide crystals that are arranged in a sphere 1.32 meters in diameter and surrounded by phototubes. The detector saw use in hadron spectroscopy at the SPEAR storage ring in Stanford. It is now in use at the DORIS II storage ring at DESY in Hamburg. (Courtesy of William W. Ash and Elliott D. Bloom, Stanford Linear Accelerator Center.) Figure 1

Hadron spectroscopy and quarks

The numerous mesons and baryons are remarkably well explained as states of simple quark "atoms" that obey the same principles as ordinary atoms.

Nathan Isgur and Gabriel Karl

Our understanding of the ultimate structure of matter has advanced greatly in the last few years. It was nearly twenty years ago that Murray Gell-Mann and George Zweig gave us the key to much of this understanding with their revolutionary proposals[1] that protons, neutrons and all other strongly interacting particles—the hadrons—are made of quarks, a theretofore unobserved kind of particle. Over the last ten years, this proposal has become firmly established even though we have not observed free quarks directly. More recent research has found that forces between quarks are extremely simple and universal—the same for all types of quarks—and that these forces explain why free quarks cannot be seen. In this article we will look at hadron spectroscopy, which has been one of the main venues for this great progress.

There is much drama associated with the long history of this field, starting with the discovery of the pion and then the strange particles, continuing with the observation of the Ω^-, the J/ψ and the Υ, and going on to this day with further investigations (figure 1) and findings (see PHYSICS TODAY, January 1975, page 17, and October 1977, page 17.) However, to simplify the discussion, we will not take a historical approach, even though this risks giving the misleading impression that the field is just as much the product of theory as experiment. The reader should see our adopting this approach simply as an indication that the field has now reached a high level of maturity and understanding. Unfortunately, this approach also fails to give proper credit to many pioneers of the quark model—for example, Richard Dalitz, O. W. Greenberg, Harry Lipkin and Giacomo Morpurgo.

The hadron as an "atom." We know from atomic physics that spectroscopy can reveal not only the constituents of a system, but also the nature of the forces between these constituents (figure 2). Although we have a theory of hadrons, in discussing hadron spectroscopy we will rely very much on an analogy between quark "atoms"—the hadrons—and ordinary atoms. Consequently, after we make a few comments on the theory of hadrons, we will give a quick description of the spectrum of the simple "atom" positronium, made of an electron and a positron. We then discuss the spectrum of the analogous system "quarkonium," consisting of a very heavy quark q and its antiquark $\bar{q}$. We relate this heavy $q\bar{q}$ system to the known $c\bar{c}$ and $b\bar{b}$ systems (for a list of quark names, see the table on page 40), and eventually, by extension, to quark–antiquark systems made of light quarks, corresponding to the ordinary mesons. The other simple kinds of quark atoms, called baryons, have no immediate atomic analog because they are made of three quarks, but these qqq systems are the most familiar hadrons, because ordinary protons and neutrons are baryons. Our discussion of quark atoms concludes with a quick look at this interesting but more complicated system.

At the end of the article, we discuss some current issues in hadron spectroscopy on which there is yet no consensus. One such topic is the question of the existence of other classes of hadrons besides $q\bar{q}$ and qqq, including some that contain no quarks at all. Finally, we sketch briefly some of the avenues of theoretical research that are being explored in an attempt to put our understanding of hadrons on a more rigorous footing.

Quantum chromodynamics. The basic theory of hadrons, quantum chromodynamics,[2] is not yet fully understood. For example, despite much effort, it is not certain that the absence of free quarks—called quark confinement—is predicted by QCD, although it now seems likely to be. While we believe that QCD is the correct theory underlying hadrons, as already mentioned, we rely here whenever possible on the analogy between quark atoms and ordinary atoms to allow the readers to use their familiarity with atomic physics. Because the analogy is imperfect, this is not always possible, so we shall begin by saying a few words about QCD.

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Quantum chromodynamics is a simple generalization of quantum electrodynamics—the theory of photons and electrons that underlies atomic physics. Instead of the electron, QCD has the quark; instead of the photon it has the gluon. The generalization of QED is that while electrons come in only one (electric) charge "state," quarks come in three ("color") charge states. Generalizations of this class were invented by C.-N. Yang and Robert Laurence Mills in 1954 for other reasons; they are called non-Abelian gauge theories. As we will see, use of the term "color" is based on an analogy with ordinary color, although the color state of a quark has nothing to do with the colors of ordinary light.

Emission or absorption of a photon leaves an electron as an electron, but emission or absorption of a gluon can cause a change of the color state of a quark. Thus, a crucial consequence of the generalization of QED is that gluons, unlike photons, must carry color charge so that it can be conserved in such transitions. Despite this difference, the resulting correspondence between QED and QCD is very deep; for example, we shall refer below to the color-electric fields and color-magnetic fields of QCD. In fact, in circumstances where we can ignore the gluon color charge, the analogy is nearly perfect, and that is why ordinary atoms are a good analog to quark atoms.

The table on page 40 lists the kinds of quarks, with their masses and charges. One should not confuse the multiplicity in the table—the "flavors"—with the three colors required by QCD: Each of the six quark flavors comes in the three colors. The table also lists the other kind of fundamental particle—the leptons—both for completeness and to illustrate the close connection apparent between these two kinds of particles, a connection that is the basis for much current speculation.

Positronium

The atom that is most useful for our analogy is a bound system of an electron and its antiparticle, the positron. This system is very similar to the hydrogen atom. In the approximation where we ignore small effects due to spins and magnetic forces, the main difference is that in hydrogen there is a great asymmetry in the constituent masses, so that the electron orbits a proton that is almost at rest. In positronium, the electron and positron orbit symmetrically about the point midway between them, a motion that is kinematically equivalent to that of an electron of mass $m_e/2$ moving about a fixed proton. The attractive force, due ultimately to "virtual photon" exchange, is just a Coulomb's law force. It leads to a ground state at energy $-\frac{1}{4}m_e\alpha^2$ (an S state, which carries no orbital angular momentum), two excited states at $-\frac{1}{16}m_e\alpha^2$ (one an excited S state, 2S, and the other a P-state, 1P, which carries one unit of orbital angular momentum) followed by three states (3S, 2P, 1D) and so on, as figure 3 shows. Here α is the fine-structure constant, which is approximately $\frac{1}{137}$.

The degeneracy of the 2S and 1P levels, is an example of a peculiarity of the Coulomb force; this equality is broken when various subtle effects are taken into consideration. There are also important corrections that split the individual levels. Foremost among these corrections are various magnetic effects associated with the intrinsic spins of the electron and positron. In the ground state, for example, the total angular momentum J is due entirely to the intrinsic spin S, because the orbital angular momentum L is zero, so the total angular momentum is 0 for antiparallel spins, and 1 for parallel spins. There are thus really two ground states, which in the spectroscopic notation $n\,{}^{2S+1}L_J$ are $1\,{}^1S_0$ and $1\,{}^3S_1$, $n = 1$ denoting the first S state. These two states are slightly split in energy, primarily by the interaction of the magnetic dipoles of the electron and positron, with the "singlet" 1^1S_0 lower in energy than the "triplet" 1^3S_1.

This splitting pattern is repeated for excited S states, but matters become more complicated if the orbital angular momentum is not zero, because then the magnetic moments of the electron and the positron interact not only with each other, but also with the magnetic fields produced by their orbital motion. These same kinds of splittings occur in the hydrogen atom, where magnetic effects involving the electron magnetic moment give rise to the so-called "fine structure," and those involving the much smaller proton magnetic moment account for the "hyperfine structure." In positronium this distinction is obviously not useful, and it is simpler to refer to spin–spin and spin–orbit effects. The result, shown in figure 3, is that a state such as 1 P is split into a characteristic pattern of four levels, $1\,{}^1P_1$, $1\,{}^3P_2$, $1\,{}^3P_1$ and $1\,{}^3P_0$. Aside from these spin-dependent deviations from the pure Coulomb spectrum, there are various other relativistic corrections that play an important role in high precision comparisons to experiment. In addition, there are some subtle effects that arise in QED as corrections to the $1/r$ potential that produce small splittings between, for example, the otherwise equal 1 S and 2 P states. This briefly summarizes the physics of positronium, where experiment agrees with theory to high precision.[3]

The lifetimes of the two states $1\,{}^1S_0$ and $1\,{}^3S_1$ are also of use in the analogy to quarkonium physics. No state of positronium lives forever. These two states decay by annihilation of the e^+e^- pair into two and three photons with lifetimes of about 10^{-10} sec and 10^{-7} sec, respectively, as predicted by QED. That ${}^1S_0$ decays to two photons and ${}^3S_1$ decays to three photons is a consequence of their internal structure and is usually summarized by saying that the ${}^1S_0$ and ${}^3S_1$ wavefunctions are symmetric and antisymmetric, respectively, under "charge conjugation" of matter into antimatter. The requirement that an extra photon be emitted suppresses the rate for ${}^3S_1$ decay by a factor of order α.

Quarkonium

After this brief review of positronium, we are in a position to understand quarkonium, a bound system of a quark and an antiquark. Whereas the forces determining the positronium spectrum are electromagnetic, in quarkonium they are "color-electric" and "color-magnetic." Even though the general behavior of these forces is not yet known, *at short distances* they are very analogous in form to electromagnetic interactions. Thomas Appelquist and H. David Politzer noticed[4] that consequently a heavy $q\bar{q}$ system, which would have a very small radius at low excitation, will feel only a Coulomb-like force and be analogous to positronium.

Figure 4 shows the low-lying spectrum of a hypothetical $q\bar{q}$ system with a quark mass that is a million times greater than the electron mass. The larger mass and the fact that the appropriate strong coupling factor α_s is greater than the fine-structure constant α lead to some scale changes, but a comparison to figure 3 shows the validity of the analogy to positronium.

Unfortunately, there is no known quark for which the analogy is as perfect as this. To understand fully the properties of the known quarkoniums we must also consider the breakdown of this analogy at large $q\bar{q}$ separation. Figure 5 summarizes the situation by showing how the interquark force varies as the potential begins to deviate from Coulombic $(1/r)$ behaviour at distances on the order of the radius of the proton, and becomes a linear potential, corresponding to a constant force, at larger distances between the quark and the antiquark. This transition leads to the confinement of quarks and explains why we don't see free quarks in nature. This feature of the quarks and QCD is one of the most interesting of the hadronic world, and we will comment on it later.

With this confinement force in effect

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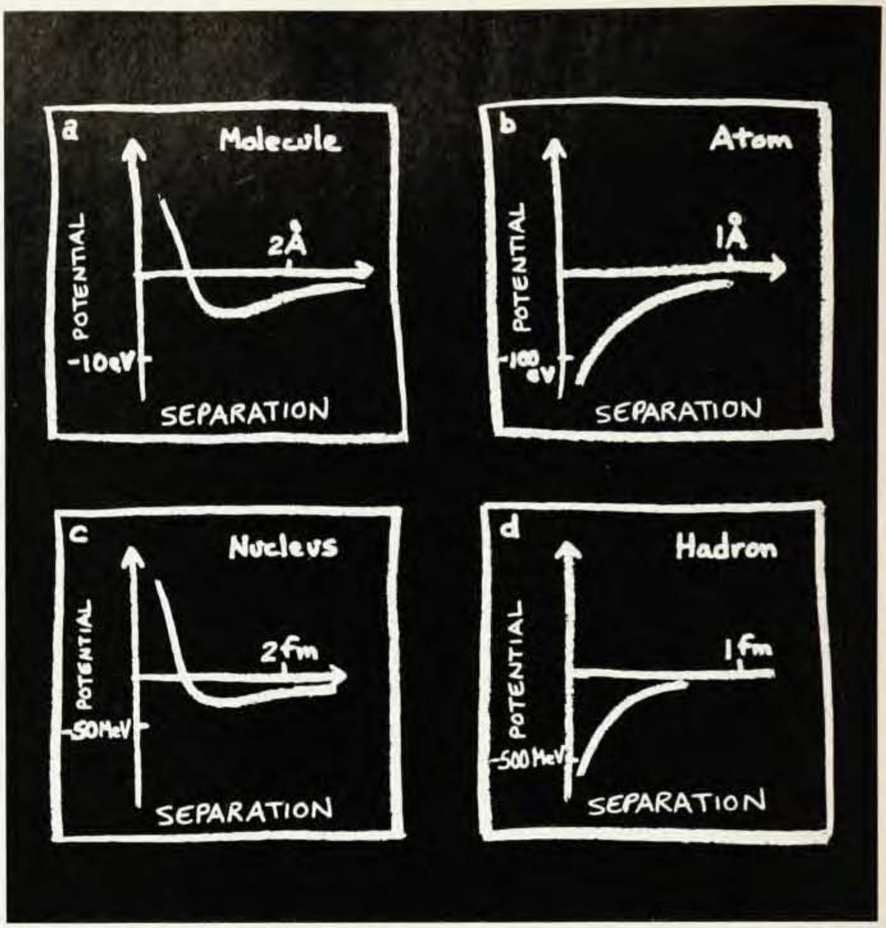
along with the short-range Coulomb-like force, a system large enough to feel the region beyond 10^{-16} meters will exhibit characteristic deviations from the Coulomb spectrum. For example, a 1 P level would appear at an energy well below that of a 2 S level. Such deviations are in fact observed in the spectra of $c\bar{c}$ and $b\bar{b}$, which we discuss next.

Real quarkoniums. We can immediately see some of the consequences of the force curve (figure 5) for real quarkoniums[5] by referring to figure 6, which shows the spectra of the $b\bar{b}$ and $c\bar{c}$ systems. The coincident discovery in 1974 at Brookhaven National Laboratory[6] and the Stanford Linear Accelerator Center[7] of the J/ψ particle, which would prove to be the $1\,^3S_1$ state of $c\bar{c}$, created a sensation in particle physics and certainly led the way to the general acceptance of quarks as the constituents of hadrons. The heavier $b\bar{b}$ system was discovered at Fermilab[8] in 1977, once again through its $^3S_1$ states. Although at first glance the $b\bar{b}$ and $c\bar{c}$ spectra look very different from the Coulombic $q\bar{q}$ spectrum of figure 4, by adding the effects of the linear potential and choosing the appropriate quark masses, one can transform this hypothetical $q\bar{q}$ spectrum into the observed $b\bar{b}$ and $c\bar{c}$ spectra. Because these systems are very clean experimentally, with many narrow states accessible simultaneously, comparison of measurements and figure 3 makes it immediately apparent that one is dealing with a spin-$\frac{1}{2}$ fermion–antifermion system, just as in positronium. Everyone agrees, these are the spectra of composite systems.

The most detailed studies of these two systems[9,10] have been carried out at electron–positron storage rings—with SPEAR in Stanford, DORIS in Hamburg, CESR at Cornell (see PHYSICS TODAY, October 1980, page 19), VEPP4 in Novosibirsk and DCI at Orsay. In these machines, countercirculating electrons and positrons annihilate, to rematerialize occasionally as $q\bar{q}$ systems. Because the annihilation occurs mainly into a state with the $^3S_1$ quantum numbers, as the collision energy is varied, one can detect the sequence of $n\,^3S_1$ states as peaks in the e^+e^- annihilation cross section. One can observe states with other quantum numbers by studying the decay products of the readily produced $^3S_1$ states. In analogy with positronium, these states have been seen via the photon cascades shown in figure 6. It is a further test of our understanding that the rates of these radiative decay processes agree with theoretical expectations. Such measurements, we should emphasize, also serve to confirm the peculiar fractional electric charges of the quarks.



Potentials between the constituents of some physical systems. In systems that are themselves composed of composite particles, the force law is usually quite complicated in form. The molecule (**a**), composed of atoms, and the nucleus (**c**), composed of protons and neutrons, are examples. In systems composed of elementary particles, the force law can be simple, as in the atom (**b**) and hadron (**d**). Figure 2

One of the most striking features of the lowest lying $^3S_1$ levels of these heavy quarkoniums are their long lifetimes. The J/ψ ($c\bar{c}$ $1\,^3S_1$) and Υ ($b\bar{b}$ $1\,^3S_1$) each have a width of about 50 keV. If these particles were oscillating circuits, their Q values would be of order 10^6! On the other hand, typical hadronic widths are at least a thousand times larger; they are characteristic of mesons that decay through the creation of a light quark pair:

$$q_1\bar{q}_2 \rightarrow q_1(\bar{q}q)\bar{q}_2 \rightarrow (q_1\bar{q}) + (q\bar{q}_2)$$

The J/ψ and Υ are too light to do this (the mass of the decay products would be greater than the mass of the parents) and so they must decay through a much slower annihilation process[4] in which the initial heavy quarks disappear altogether:

$$q\bar{q} \rightarrow \text{gluons} \rightarrow \text{light hadrons}$$

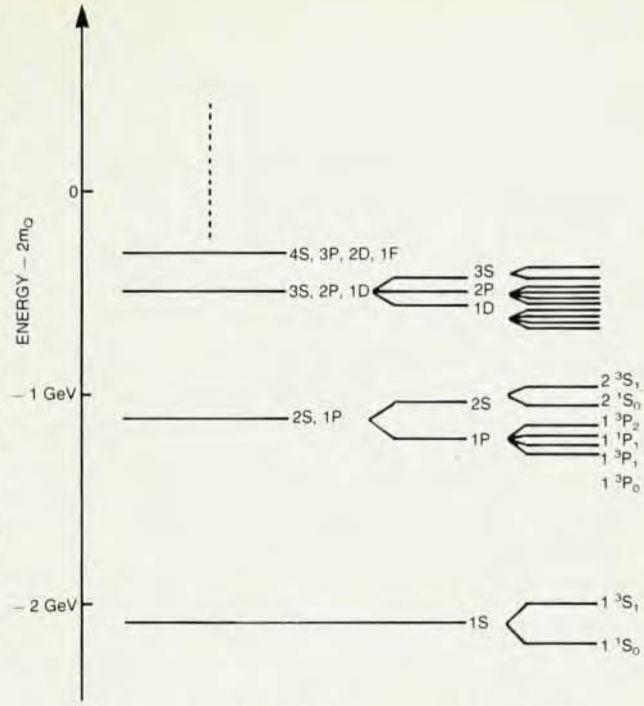
We can calculate the widths appropriate to *this* mechanism of decay of heavy quarkoniums by means of a formula analogous to that giving the width for the decay of the $^3S_1$ state in positronium. This is because at short distances, below about 10^{-15} meters, the process $q\bar{q} \rightarrow$ gluons is just like $e^+e^- \rightarrow$ photons, even though at larger distances, the gluons, unlike the photons, become confined.

Light mesons

We can take the same ideas that lead to such a successful description of the heavy quarkonium systems and apply them to mesons made of light quarks. Alvaro de Rújula, Howard Georgi and Sheldon Glashow suggested[11] this important extension in a very influential paper in 1975. Historically, the quark model emerged from the description of the light mesons and other light quark states. From the three light-quark flavors, u, d and s, one can construct 9 mesons of each space and spin configuration. By 1964, when the quark model was invented, the nine $1\,^1S_0$ states π^+, π^0, π^-, K^0, K^+, $\bar{K}^0$, K^-, η and η' were known and so were the nine $1\,^3S_1$ states ρ^+, ρ^0, ρ^-, K^{*0}, K^{*+}, $\bar{K}^{*0}$, K^{*-}, ω and ϕ. Although the character of these states precluded the clear demonstration of compositeness that became possible later with the discovery of heavier quarkoniums, and even though a very precise description is not yet available, light quarkoniums are quite similar in character to their heavier cousins $b\bar{b}$ and $c\bar{c}$. It is actually rather hard to avoid this conclusion. As figures 7 and 8 show,[12] the basic character of the spectra is unchanged from what we saw in $c\bar{c}$ and $b\bar{b}$, even though the non-Coulomb, spin–spin and spin–orbit effects are relatively larger.

Notice in particular that the lightest charged non-strange mesons, the π and the ρ, are analogs of the ground states $1\,{}^1S_0$ and $1\,{}^3S_1$ of the hydrogen atom: The ρ–π splitting that makes the π the lightest hadron is the analog of the 21-cm hydrogen line of radioastronomy fame. As with heavier quarkoniums, the internal structure of the states, deduced by dynamical properties such as radiative transitions, is consistent with this picture.[13]

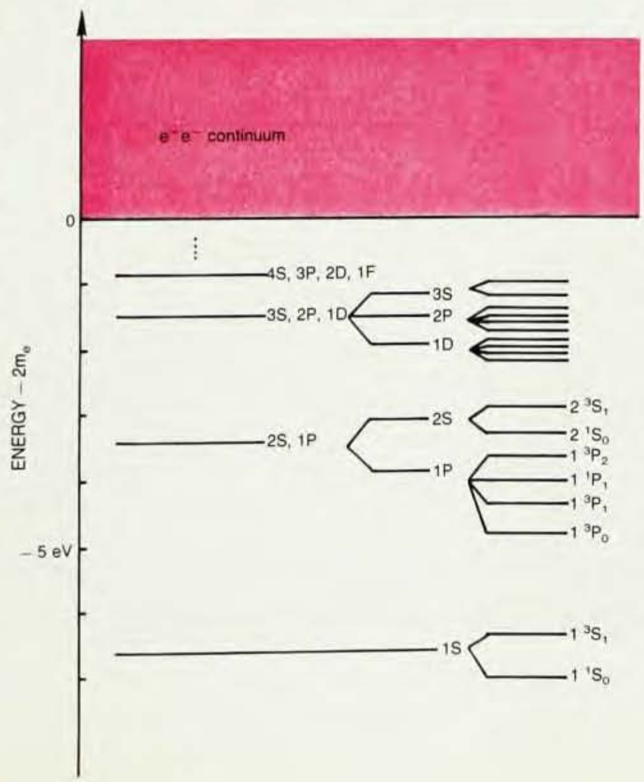
We can now see that the complexity of meson spectroscopy is mainly illusory. In a world with six quark flavors, mesons come in 36 varieties (six quarks $\times$ six antiquarks), each of which can be in the various orbital and spin states we have discussed, so it is not surprising that hundreds of such states are known. We hope that the reader now appreciates that they are all examples of the same underlying principles, similar to those underlying the hydrogen atom itself.



Heavy quarkonium. This diagram shows the low-lying energy levels in a hypothetical heavy quark–antiquark system. As in positronium (figure 3), the energy levels are determined mainly by a Coulomb-like force. The levels on the left are shown to scale. Their splittings are shown very schematically, multiplied by a factor of roughly ten. As in figure 3, the ordering of states within an orbital is correct, but those between orbitals are distorted for simplicity of presentation. The "Bohr radius" of this $q\bar{q}$ system is about 7.7×10^{-18} meters. Figure 4

Light baryons

The hadrons that have been known the longest are the baryons. Exemplified by the proton and neutron, these particles are composites of three quarks. One might be puzzled, based on what we said above, that three quarks are bound at all. The binding is intimately connected with the three-dimensional (non-Abelian) nature of color charge. It would take us too far afield to explain this interesting fact, but a crude analogy to ordinary color may be useful. By taking two quarks with different primary colors one can produce the complement to the third primary color. Thus, a given quark would see the other two quarks as having effectively the color of its antiquark, and so will bind to them.



Positronium energy levels. In the e^+e^- "atom," the energy levels are determined mainly by Coulomb's law. These are shown on the left to scale. The splittings of these Coulomb levels appear very schematically on the right, multiplied by a factor of roughly a thousand. The ordering of states within an orbital is correct, but those between orbitals are distorted for simplicity of presentation. The "Bohr radius" of positronium is 1.06×10^{-10}m. Figure 3

While confinement is still not understood in detail, it is easy to construct models that incorporate its effects. One can then compare the expectations of these models with the spectra of baryons.[14] The experimental informatiion that exists on baryons is generally of very high quality, but it is almost exclusively restricted to baryons made of the three light quarks u, d and s. There are, unfortunately, no known baryons with the simplicity of heavy quarkoniums, so our discussion will focus on light baryons.

It is not as easy to enumerate the possible states of baryons as it is for mesons, because the Pauli principle forbids two identical quarks to be in the same state. To circumvent this difficulty for the moment, we first discuss baryons composed of the three distinct quarks, u, d and s, so that the Pauli principle is irrelevant. Consider the u and d quark in such a baryon. They can, as in a meson, have an orbital angular momentum L_{ud} of 0,1,2, . . . and a spin S_{ud} of 0 or 1. Next consider the motion of the ud pair as a unit relative to the s quark: Add the spin of the s quark to the spin S_{ud}, and add the orbital angular momentum $L_{(ud)\text{-}s}$—which is 0,1,2, . . .—to L_{ud}. If the spin S_{ud} is 0, then the total spin S_{uds} is $\frac{1}{2}$; if the spin S_{ud} is 1, then the total spin S_{uds} can be $\frac{1}{2}$ or $\frac{3}{2}$. In the ground orbital state we expect the orbital angular momenta L_{ud} and $L_{(ud)\text{-}s}$ both to be 0, so there will actually be three possible "ground" states (corresponding to the three possible spin states): two with total angular momentum $J=\frac{1}{2}$, and one with $J=\frac{3}{2}$.

As in the case of mesons, spin–spin interactions via color magnetism will produce splittings between these states. It can be shown that the lowest-lying state will have spin $S_{ud} = 0$ and total angular momentum $J = 1/2$ (it is the $\Lambda(1115)$), that the next highest state will be the other $J = 1/2$ state and will be about 80 MeV higher in mass (the $\Sigma^0(1195)$), and that the $J = 3/2$ state (the $\Sigma^{*0}(1385)$) will be yet another 190 MeV above the Σ^0.

As with mesons, there are excited baryons corresponding to orbitally excited quark states. If either orbital angular momentum L_{ud} or $L_{(ud)-s}$ is excited from 0 to 1, the results are states that lie roughly 500 MeV above the ground state. Experiments have in fact observed these "negative parity" baryons. More complicated states result if the orbital angular momentum is excited by two units. This gives rise to the "positive parity" excited baryons, which have also been observed. In each case these orbital states have to be coupled with the three spin states described above, so a very rich spectrum emerges.

If two or more quarks are identical, then the exclusion principle eliminates certain states that would otherwise have been allowed. Precisely which states are ruled out is intimately connected to color. We mentioned earlier that the three quarks in a baryon are in a state where each quark has a different color; QCD in fact requires the quarks to be in a state that is totally antisymmetric in their colors. Consider, for example, the ground state of uuu. With both internal angular momenta zero, the exclusion principle requires the system to be in the *symmetric* $S = 3/2$ state, in which the three spins are aligned, and excludes both $S = 1/2$ states. Indeed, the only low-lying doubly charged baryon is the Δ^{++} at 1235 MeV, which has a total angular momentum J of $3/2$.

It turns out that in the case of uud and ddu—appropriate to the proton and neutron, respectively—both the spin-$3/2$ state and one of the spin-$1/2$ states are allowed. The forbidden spin-$1/2$ state has $S_{uu} = 0$. One can show that the color-magnetic interactions in this case lower the energy of the spin-$1/2$ state, explaining why the proton and neutron are the lightest baryons. Their splitting from the four spin-$3/2$ states, Δ^{++}, Δ^+, Δ^0 and Δ^- is, like the ρ–π splitting, a spin–spin interaction. It is interesting to note that the four $S = 3/2$ states have nearly the same mass. This fact is connected to the near equality of the u and d quark masses, and is described formally as an "isospin" symmetry.

Experimentally,[12] the rich spectroscopy and dynamics of the non-strange baryons uuu, uud, ddu and ddd are known quite well up to masses of about 2 GeV, and the observed states seem to be in good correspondence with those one expects. As with the light mesons, the uncertainties are such that one's expectations are less quantitative than in a heavy quark system. Nevertheless, there seems to be no good reason to doubt that in light baryons, as in light mesons, we understand the main principles of quark atoms.

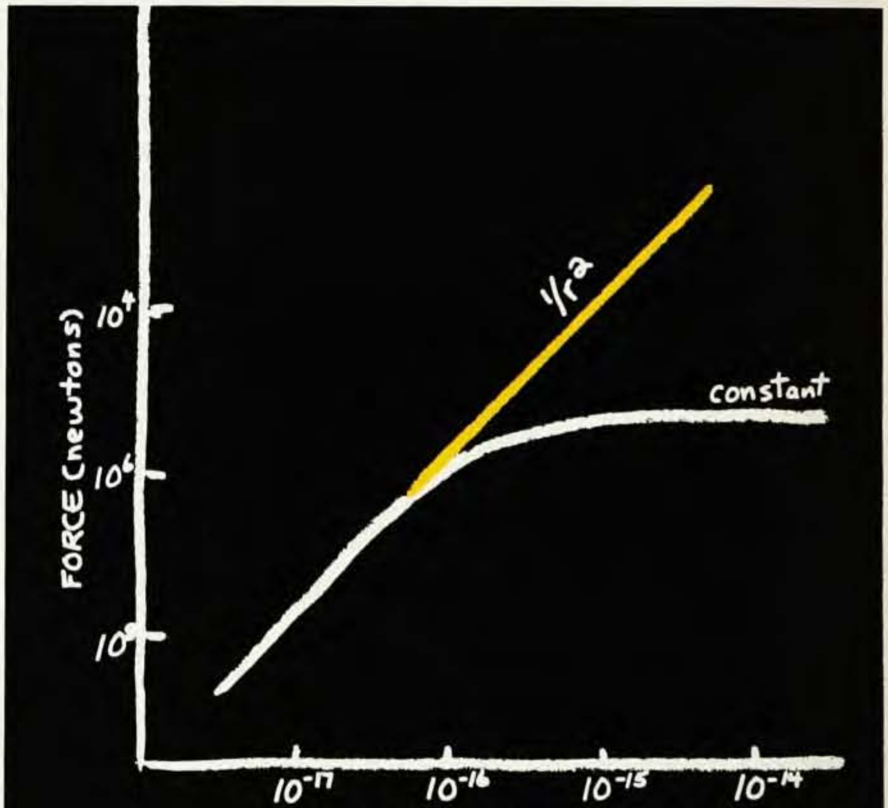
Fundamental particles

	Symbol	Mass* (GeV)	Charge
Quark flavors			
Down	d	~0.35	$-1/3$
Up	u	~0.35	$+2/3$
Strange	s	~0.55	$-1/3$
Charm	c	~1.7	$+2/3$
Beauty (or bottom)	b	~5.2	$-1/3$
Truth (or top)	t	hypothetical	$+2/3$
Lepton flavors		(MeV)	
Electron	e	0.511	-1
Electron neutrino	ν_e	$<5\times10^{-5}$	0
Muon	μ	106	-1
Muon neutrino	ν_μ	<0.5	0
Tauon	τ	1784	-1
Tauon neutrino	ν_τ	<250	0

*For quarks we quote the "constituent" quark masses.

Other kinds of hadrons. Once the idea of a quark substructure for the known hadrons is established, it is natural to ask whether $q\bar{q}$ and qqq are the only configurations allowed. Strictly speaking, we know they are not, because nuclei exist: They are quark molecules, constructed mainly from quark atoms—the protons and neutrons. Another possible exception, suggested by MIT particle physicist Robert L. Jaffe, is that the two puzzling states S* and δ observed just below the $K\bar{K}$ threshold may be $qq\bar{q}\bar{q}$ states. If so, it is plausible that they are meson molecules. There may be some other non-molecular "multiquark" states with some degree of stability, but the theory of such states is not well understood and the subject remains contoversial. In fact, of the many hundreds of hadrons now known, there are no established states other than nuclei that do not fit into either the $q\bar{q}$ or qqq family.

There are, nevertheless, strong theoretical indications for other kinds of hadronic matter, and we would like to give a quick survey of the kinds of new objects that people are seeking and debating, without stressing our own prejudices, which are changing rapidly. In the preceeding discussion we have



Force between quarks as a function of separation. At distances between 10^{-16} and 10^{-15} meters, the interquark force becomes confining with a constant value of about 1.6×10^5 newtons. This universal potential governs all quarkoniums, ($b\bar{b}$, $c\bar{c}$ and so on), which differ only in the masses of their quarks. Note that the scales on both axes are logarithmic so that the Coulomb force becomes a straight line. Figure 5 only briefly mentioned gluons because, as is the case for photons in ordinary atoms, their presence is not manifest in mesons and baryons. However, just as a world without charged matter would still have a spectrum of photons, one without quarks would still have a spectrum of gluons. In the latter case, however, we expect the spectrum to be very rich in structure because QCD predicts confinement of the gluons. This confinement has led to the widespread belief that hadrons of pure glue should exist.[15] (See PHYSICS TODAY, July 1981, page 20.) Such states are predicted in many models in which they have masses similar to those of quark atoms. These models also indicate the existence of "hybrid" states, in which both quark and gluonic degrees of freedom are excited simultaneously. If these expectations are borne out by experiment, then hadron spectroscopy has not only a very rich, but also a very eventful, future.

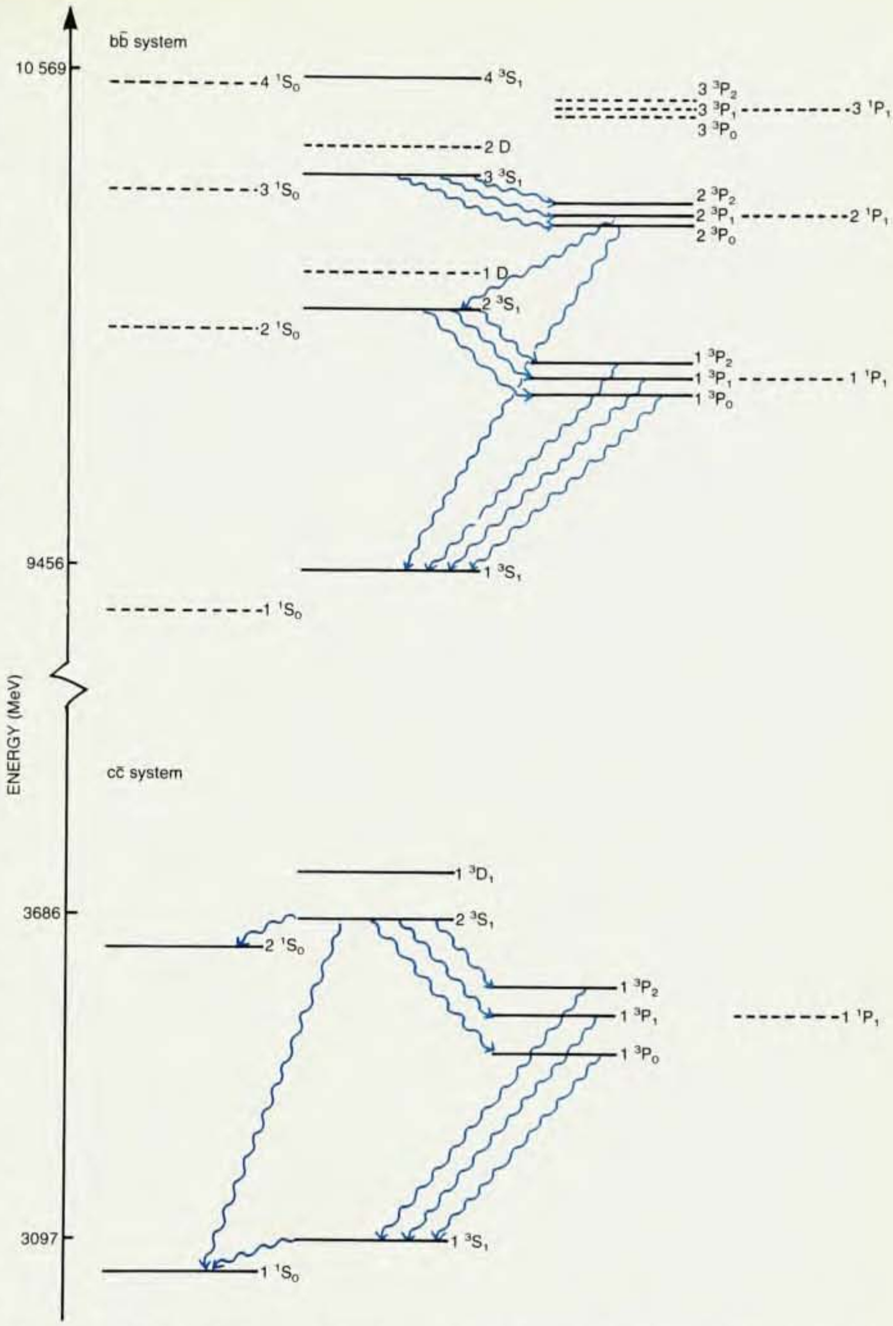
QCD and phenomenological models

While not our main topic, the theory underlying hadron spectroscopy is obviously vital to this subject. As we have implied, the consensus is that the correct theory is quantum chromodynamics. However, QCD is still in its infancy, with rather few rigorous results having been derived from it so far. In contrast to quarks and gluons, electrons and photons are only weakly coupled, so that a perturbative treatment is natural and successful. With quarks and gluons this is only the case at short distances. However, in the confinement region most relevant to light hadrons, the coupling between quarks and gluons becomes very strong, perturbative methods are bound to fail, little is understood and even less is rigorously proved. Research is showing some progress, however, along many lines of attack.

One method for understanding the strongly interacting regime of QCD was pioneered in the Soviet Union. This technique uses our understanding of interactions at short distances, together with causality, conservation of probability and other general principles, to extract information about hadrons.[16]

Other progress comes from numerical work on field theories. By putting fields on a physically fictitious lattice[17] in space and time, one can apply many techniques from solid-state physics. This method, known as "lattice gauge theory," has already given indications that at large distances quarks are confined in QCD with a linear potential. However, so far these methods have produced only semi-quantitative results in spectroscopy.

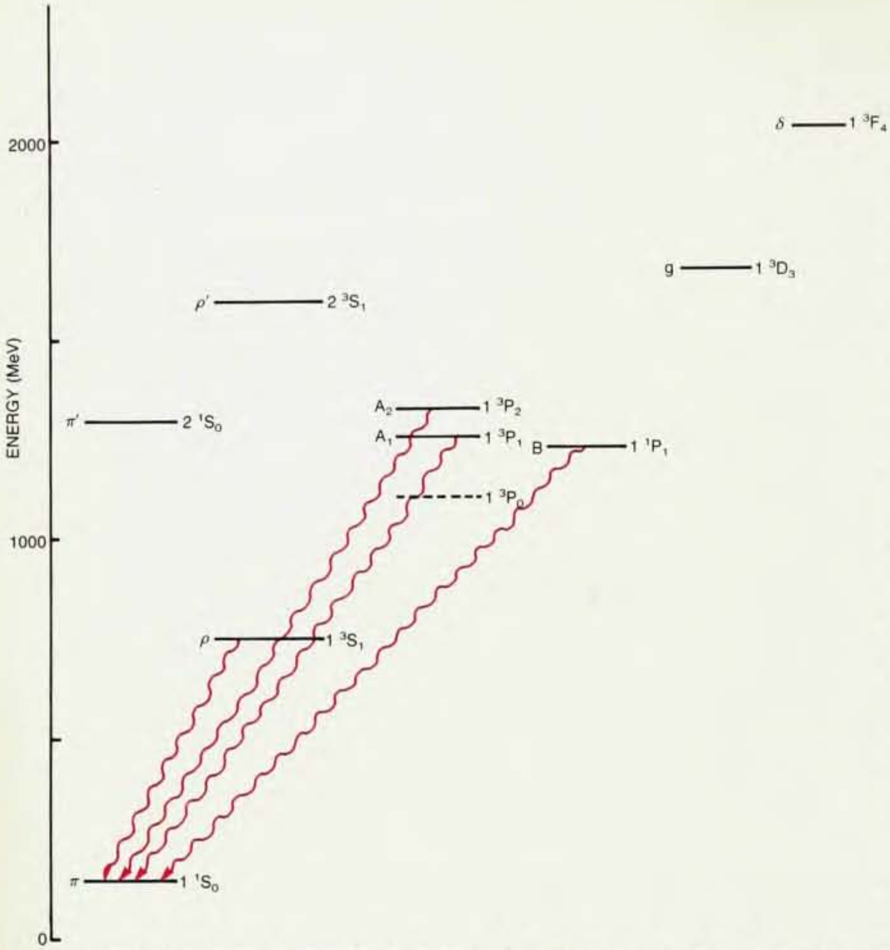
Lattice gauge theory may eventually lead to accurate approximations of QCD, but for the foreseeable future we



Spectra of the $b\bar{b}$ and $c\bar{c}$ systems. These energy-level diagrams show the low-lying spectra of the two known heavy quarkoniums. The solid lines are observed states, dashed lines are expected ones and wavy lines are observed photon transitions. Figure 6

will continue to need models to study the known hadrons. One possibility is to find simplified quantum field theory models that can approximate QCD. (With this possibility in mind, Gerard 't Hooft and Edward Witten have noted that as the number of colors goes to infinity, the theory simplifies greatly.) The most likely source of progress, however, is the continued use and refinement of the rather naive atomic models of the type we have been discussing.

Some models for confinement are not based on the concept of a potential. Foremost among these is the "bag model"[18] based on field theory in a cavity, or "bag," that confines quark and gluon fields. The physical idea behind this model is that space—the vacuum—has two phases that can coexist: one in which quarks and other colored objects can live, and the other from which they are expelled. This forms the bag, within which quarks can move relativistically, interacting with each other via short-range QCD forces. The resulting model gives a successful description of ground-state hadrons and their many properties. Pure glue and hybrid states have a natural place in this model. One can excite gluon modes inside the bag much as one can excite microwaves in a cavity. In spite



Light mesons. The diagram shows some low-lying $u\bar{d}$ mesons. The solid lines are observed states, the dashed line is an expected state and the wavy lines are observed photon transitions. Figure 7

of the obvious physical appeal and the many successes of this model, its spectroscopic usefulness has so far been limited. This is parimarily due to technical difficulties with the motion of the center of mass in excited states, difficulties that are common in relativistic models.

Another phenomenological model of confinement is the "flux-tube" or "string" model. In this model the dipole color-electric field around a $q\bar{q}$ system collapses into a tube or string of flux that ties the quark to the antiquark. At large distances this produces a constant force between the quark and antiquark: the "tension in the string." In baryons, the flux collapses into a Y shape with the three quarks at the ends. In this model, one identifies pure glue states with closed loops of color-electric flux, and hybrids with states in which the string between quarks is excited transversally.

The most common of the naive confinement models, on which we relied throughout this article, are those based on a potential, which we relied on throughout this article. As we have emphasized, these are particularly apt for heavy quark–antiquark systems but have shown their utility throughout hadron spectroscopy. One merit of nonrelativistic models is that they allow the separation of the relative motion of quarks from the overall motion of the hadron. It is our prejudice that in spectroscopy it is better to have the right number of moving parts, moving at the wrong speed, than to have the wrong number of moving parts moving at the right speed.

All these models are, to some extent, complementary. For example, work with flux-tube and bag models has yielded suggestions for potentials. In any event, until more fundamental studies are successful, it is surely those features common to all models that we should emphasize.

Loking back over the last ten years, it is clear that hadron spectroscopy has made great progress. Quantum chromodynamics has given us a unified picture of quarks and the forces between them, and has provided a broad and detailed understanding of vast amounts of experimental data. Nevertheless, our knowledge of hadrons is far from complete both quantitatively and qualitatively. On the quantitative side, there is hope of progress in extracting from QCD predictions of a precision that rivals that of atomic physics. Even more exciting is the prospect of discovering new types of hadronic phenomena that have eluded us so far. While relative to the new frontiers of high-energy physics, hadron spectroscopy is now *terra cognita*, there is still much to explore.

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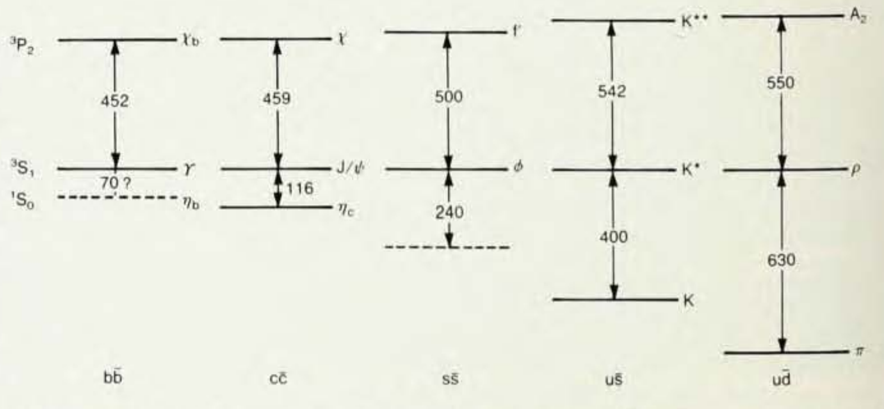

Transition from heavy quarkoniums (left) to light quarkoniums (right), showing the evolution of the $^1S_0$ and $^3P_2$ levels relative to the $^3S_1$ level. The η_b is not yet observed, so its expected value is shown. The $s\bar{s}$ $^1S_0$ state is not directly measurable; the mass shown was extracted theoretically from the η and η' masses. Figure 8

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