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A Nonstandard Model

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An elementary-particle picture developed primarily by Barut as an alternative to the standard model is re-examined. This model is formulated on the basis of strong short-range magnetic interactions among the stable particles (p, e^-, ν) *and at present is able to account qualitatively for most of the known phenomena.*

For more than three decades high-energy physicists have labored to assemble a wealth of data into a modern picture of elementary particles. The result, described principally in terms of quantum field theory (QFT), is a body of conventional wisdom generally referred to as the *standard model.* This development has not proceeded without reasoned disagreement, however, and a parallel body of other views exists, which have been less successful in reaching fruition. In particular, the work of Asim Barut and colleagues has aimed to construct a more transparent theory of particles that avoids their instantaneous appearance from, and disappearance into, the vacuum—and which abjures the introduction of new forces in a picture already unified by the electromagnetic field. Guided by a penchant for simplicity and a vision of "the way it ought to be", Barut has provided the skeletal structure for an alternative picture that actually turns out to have a measure of flesh on it. The time seems quite appropriate to gather together and review the essential features of these alternate views, due primarily to one who is very definitely a nonstandard model himself.

1. THE STANDARD MODEL

Regardless of what the future holds, the so-called standard model of particles and fields, consisting of the electroweak theory and quantum

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chromodynamics (QCD), can count a number of impressive successes. Among these are predictions of the massive vector bosons, weak neutral currents, and certain parity violations manifested as asymmetries in various cross sections. In addition, the observed hadronic weak-decay rules are understood within the theory, as is the nonobservation of otherwise-expected decays (for example, $\mu^+ \to e^+\gamma$, $K_L \to \mu^+\mu^-$). The abundant data on e^+e^- annihilation, and on neutrino scattering, provide further support for the model.[1] In view of this plethora of positive features, it would seem almost churlish to sound a negative note without good cause. We leave it to the reader to render that judgment.

To begin, the Higgs boson and the top quark have yet to be seen, and both quarks and gluons have been declared "unseeable" in the wild (asymptotic freedom). In the simplest version there are some 21 adjustable parameters, including particle masses and strengths of the forces. Although the electroweak theory is constructed to exhibit maximal parity violation, this is not explanatory, but put in by hand, and results in 33 different fermion multiplets. Both color charge and family structure were crucial to the development of the quark model and electroweak unification—the former is not observed, however, and the latter, along with flavor symmetry, is not understood at all within the model. One might have expected a simple quantitative explanation of the neutron–proton mass difference within the theory, but there is difficulty in getting this even qualitatively correct.

It can be argued that these problems may eventually work themselves out, though it is also possible that the tortuous path taken to the standard model is an indication that it carries a bit too much baggage. It is a complicated theory, and much of the complexity derives from its basis in QFT, which was indeed resurrected from the dead by development of the electroweak theory.

In originating QFT, Fermi went beyond the need for a mathematical means of describing pair processes, and envisioned (e^-, $\bar{\nu}$, p) to be created spontaneously in β-decay, while the neutron vanished in the same way.[2] His stated motivation was the inadequacy of contemporary relativistic theories of leptons to account for binding of these particles at nuclear dimensions. As with quantum electrodynamics (QED), any other QFT must also be renormalized, and this was accomplished in the electroweak case by making it a gauge theory. But the gauge symmetry must be exact, so the notion of spontaneous symmetry breaking through the Higgs mechanism was introduced to conform to the observation that the symmetry is actually broken in nature.

Although the notion of symmetry has historically had great appeal for the human psyche, and has been a powerful tool in acquiring an understanding of nature, the insights achieved have almost always referred to spacetime symmetries. Extensions to internal symmetries of elementary particles are of more recent origin, and almost always refer to broken symmetries. Indeed, it is only the lower-dimensional finite subgroups that are discussed, and the full continuous Lie-group structure seems superfluous—otherwise, one would expect all subgroups to appear. Thus, one might question the need to gauge these symmetries as a fundamental requirement, or to consider the multiplet structure much more than a useful classification scheme. No particularly deep meaning has been uncovered for the associated internal quantum numbers, in any event.

For some, these twists and turns have a strong flavor of teleology, and one wonders if there might not be a simpler, more transparent way—a way that avoids second quantization altogether. (Although we have yet to understand the physical mechanism underlying pair processes, only QFT has been able to supply a mathematical description of the phenomena; for this reason it seems practicably difficult to abandon it completely until another alternative emerges.) It was already known in the late seventies that theories could be developed reproducing the data of that time without gauging or spontaneous symmetry breaking.[3,4]

Is it possible that the electromagnetic unification provided by Maxwell has not been taken far enough, and there exists an already-unified theory of elementary particle interactions without introducing extraneous forces? Ockham's razor alone, or economy of theoretical constructs, suggests that what we call weak and strong forces may be "fictitious" in the same sense that we classify Coriolis, centrifugal, and chemical forces, and that things ought to be made of those entities into which they decay. (This latter desideratum, of course, has always been in conflict with the conventional interpretation of quantum theory itself.)

An evident lepton–quark symmetry in particle physics has long been held remarkable. One implication is that it may be profitable to attempt to construct a leptonic theory of hadrons, an idea that is not particularly new *per se*.[5–7] There are obvious difficulties with such an effort, and the most serious questions to be addressed include the following:

(1) How can large hadron masses be obtained from electrons and neutrinos?

(2) How can electrons and neutrinos be bound at nuclear dimensions and localized to that size (Fermi's question)?

(3) How can the relatively large electron magnetic moment be reduced upon binding to the level of hadronic, or even, heavy-leptonic moments?

(4) How can the observed hadronic classifications be reproduced by leptons alone?

These previous attempts either introduced a new and unknown "hadronic core" that binds leptons,[5,7] or envisioned leptons as possible quark constituents.[6] The remainder of this article reviews a different approach—one that introduces no new phenomenology, but depends on the full range of electromagnetic interaction. Of necessity, the context is primarily qualitative, and the obstacles to obtaining many quantitative results at present will become apparent in the following section, though some have been booked.

2. MAGNETIC INTERACTIONS AND NEUTRINOS

Almost from the moment Pauli first perceived its existence,[8] the neutrino has had associated with it a possible intrinsic magnetic moment. Carlson and Oppenheimer[9] applied Pauli's suggested nonminimal coupling to a study of the ionization loss in interaction with electrons, and the cross section for electron–neutrino scattering via a magnetic moment ($\sim \ln E$) was re-derived by Bethe.[10] (These are, perhaps, the first known studies of neutral-current interactions!) Apparently, Pauli also suggested that the neutron may be a bound state of electron, proton, and neutrino in a deep magnetic well, but the requisite calculations were not forthcoming, and the arguments of Carlson and Oppenheimer mitigated against the model. As noted above, these notions were overshadowed by the immediate success of Fermi's weak-interaction theory.

It is curious that the electromagnetic interaction is so fundamental in nature, yet the magnetic field seems to play only a minor role as a weak perturbation in atomic physics (though often a major role in astrophysics). But at the magnetic radius $r_m \equiv \alpha \lambda_c$—about 3 fermis for an electron—the magnetic field is about 10^{13} T, and the magnetic energy outside this volume is roughly ten times the electrostatic energy.[11] Moreover, the magnetic energy of one electron in the field of another equals its rest energy at ~ 40 fermis, so that one can hardly treat magnetic interactions as small perturbations at these distances. Similar observations have been noted explicitly with respect to gravitational interactions between electrons, where it is shown that the magnetic-moment contributions dominate the field at distances between the classical electron radius and the Compton wavelength.[12,13] In addition, the importance of spin effects at high energies has become evident in the striking results of polarized proton–proton elastic scattering.[14,15]

The so-called solar neutrino problem has generated a renewed interest in a possible neutrino magnetic moment, of course, and numerous estimates suggest a (charitable) upper bound of $\sim 10^{-10}\mu_B$. Thus, the solar neutrino deficit and flux variation *may* be explained through induced helicity flips,[16] or via $\nu_e - \nu_\mu$ conversion[17] in the solar magnetic field. One would expect the existence of a small neutrino mass to enhance the possibility for a magnetic moment. A magnetic neutrino, with or without mass, is necessarily described by a 4-component Dirac bispinor,[18] and thus couples to matter and fields through the nonminimal Pauli coupling. Thus, a general Dirac particle in the presence of an electromagnetic field satisfies

$$\left[\gamma^\mu(i\partial_\mu - QA_\mu) + \frac{\mu}{2}\sigma^{\alpha\beta}F_{\alpha\beta}\right]\psi(x) = m\psi(x) \tag{1}$$

where $\mu = a\mu_B$ is an anomaly in terms of the Bohr magneton, and $\hbar = c = 1$. It is then possible to envision electrons, neutrinos, and their antiparticles interacting via effective potentials containing magnetic wells and barriers, as indicated qualitatively in Fig. 1. Such potentials are manifested by forms such as

$$V(r) = \frac{A^2}{r^4} + \frac{B}{r^3} + \frac{C}{r^2} \tag{2}$$

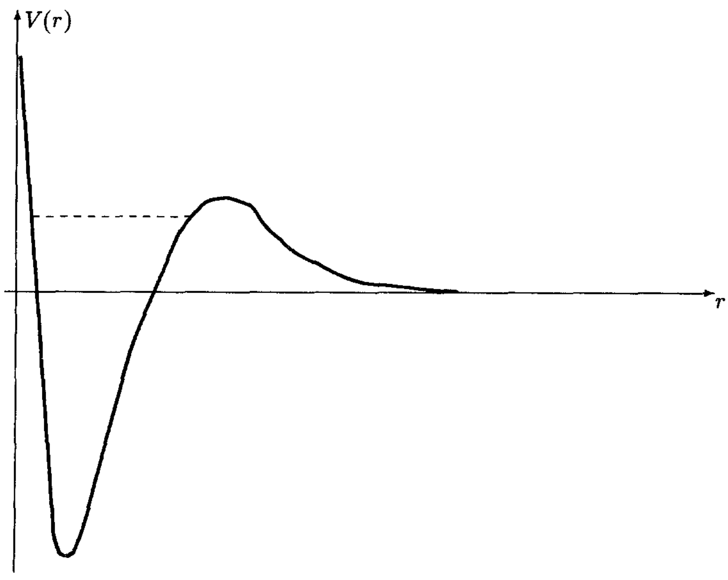

Fig. 1. Qualitative representation of the potential function $V(r)$, Eq. (2). The dashed line denotes a possible positive-energy resonance.

which has been studied by Predazzi and Regge,[19] as well as by Barut and collaborators,[20,21] and these possess some notable features. The possibility of positive-energy resonances and bound states is apparent, along with processes facilitated by resonance penetration and α-decay-like tunneling. If calculations with localized states are carried out nonperturbatively, it is possible that such potentials might actually sum certain perturbation series and render parts of the renormalization program unnnecessary. We shall see below how these possibilities might be manifested in particle processes.

Continued absence of electron structure in high-energy cross sections —down to $\sim 10^{-16}$ cm—indicates that magnetic interactions may be very important at very short distances. Indeed, just such a mechanism has been suggested as a possible source of anomalous positron peaks in heavy-ion collisions[22]—it is occasionally thought of as a "phase transition" in QED. The existence of e^+e^- magnetic resonances in terms of the anomalous magnetic moments has been studied in some detail by Barut and co-workers,[23,24] and these states are referred to as *superpositronium*. Although detailed nonperturbative calculations of the spectrum have yet to be made, possibilities abound: "The qualitative resemblance between the ψ family spectrum and positronium is certainly striking."[25]

The suggestion that charged particles possessing magnetic moments may form relativistic bound states at nuclear distances is not at all new, and was considered over 30 years ago.[26,27] These calculations were carried out classically, then subjected to Bohr quantization. The results were moderately encouraging, but were pursued no further. It is our contention here, along with Barut, that magnetic interactions among the stable fermions—e^+e^-, $e\nu$, $\nu\bar{\nu}$—might form the basis for a sound and transparent model of elementary particles. Difficulties with pair processes aside—which remain unexplained in any theory—one envisions the Dirac theory extended electromagnetically to the complete energy spectrum:

Coulomb:	$r \sim \lambda\!\!\!^{-}_c/\alpha(\sim 10^5\text{f})$	$E \sim \alpha^2 mc^2(\sim 10\text{ eV})$
Nuclear:	$r \sim \alpha\lambda\!\!\!^{-}_c(\sim 3\text{f})$	$E \sim mc^2/\alpha(\sim 70\text{ MeV})$
Supernuclear:	$r \sim \alpha^2\lambda\!\!\!^{-}_c(\sim 10^{-2}\text{f})$	$E \sim mc^2/\alpha^2(\sim 10\text{ GeV})$

Given this, at least for the sake of the argument, let us see where it leads.

3. AN ALREADY-UNIFIED PICTURE OF ELEMENTARY PARTICLES

The conventional view of the structure of matter has by and large been conditioned by experiment to be hierarchical *donward*. Thus, as accelerators have moved to higher and higher energy, we have been able to probe that structure to smaller and smaller distances. And as "particles" proliferate, it becomes desirable to organize the picture by looking for more fundamental constituents. Given the choice between an eventual limit and an infinite regress, we clearly must hope for the former if there is to be any chance of ever understanding that structure.

Rather than considering higher energy processes to probe toward yet finer structure, perhaps it is time to presume that limit to have been attained in the form of the known stable particles: p, e^-, ν, and their antiparticles. (We shall discuss the stability and structure of the proton presently.) As collision processes proceed to higher energies, it appears that these stable particles form various series of resonances with generally decreasing lifetimes, and then these resonant states eventually decay back into the stable particles of which they are composed. No doubt the number of particles is limitless in this direction, as long as the input energy continues to rise. With increasing energy new resonances, or "particles," are constructed from those more stable entities preceding them on the energy scale.

It has long been known that the data on which the quark model was originally based are open to other interpretations. For example, the particle spectrum can be made compatible with integrally charged quarks as well,[28] and below we shall see that by grouping leptons and baryons appropriately into $SU(3)$ multiplets one can establish a complete mathematical correspondence with the quark model. Moreover, from deep lepton–nucleon scattering data, one can infer two possible solutions for charges of the constituents:[29] one yielding nonintegral (including quark) charges, and another yielding $(+1, +1, -1)$ for protons and $(+1, -1, 0)$ for neutrons.

Finally, one of the most important sources of data at high energy is that from $e^+e^- \to$ hadrons. Much of this data is compatible with an $SU_c(3) \times U(1)$ theory with spontaneously broken color symmetry and integral charged quarks.[30] An important descriptor of these processes is the ratio of the hadronic cross section to that for $\mu^+\mu^-$ production:

$$\begin{aligned} R &= \frac{\sigma(e^+e^- \to \gamma \to \text{hadrons})}{\sigma(e^+e^- \to \gamma \to \mu^+\mu^-)} \\ &\simeq 3 \sum_f Q_f^2 \end{aligned} \tag{3}$$

where the sum of quark charges goes over all flavors produced at a given energy, and the factor of 3 is the number of colors. Below the charm threshold, $R \simeq 2$—but this is also the value obtained from the two stable charged leptons e^+, e_-: $\sum_l Q_l^2$. In a lepton model there is no statistics problem, hence no need for a color quantum number. (At higher energies one must include $\mu^{\pm}$ and $\tau^{\pm}$ as they become relatively stable on those time scales.)

Let us follow a number of Barut's earlier ideas along these lines,[31–34] and propose a picture based on the following propositions:

(1) The only stable particles are p, e^{-}, ν, and their antiparticles, and all other particles are constructed from the stable constituents into which they decay. (For the moment we do not distinguish between electron and muon neutrinos.)

(2) There exists a 4-component neutrino with intrinsic magnetic moment, which may possibly also possess mass; only in the asymptotically-free case does the Dirac equation split into 2-component equations.

(3) With the exclusion of gravity at this level, all particles interact through the electromagnetic field only, thereby providing an already-unified theory.

(4) Stable particles, and only those, can be created and annihilated in particle–antiparticle pairs; with the exception of normal quantum-mechanical exchange and rearrangement (Pauli principle), this is the only other dynamical mechanism proposed, though its precise nature remains obscure.

Maxwell's equations and the Dirac equation (including anomalous moments) provide the basic dynamical descriptions of the stable spin-$\frac{1}{2}$ fermions, although we recognize that this is incomplete for the proton, which is now known to be composite. The only fundamental boson is the massless photon, which for practical reasons we continue to view as the quantum of the electromagnetic field, whether or not field quantization is presumed—both field and particle manifestations provide useful tools in what follows.

By asserting electromagnetism to be the sole dynamical force, we also take the binding mechanism at high energies to be the dominance of the magnetic interactions at short ranges, as discussed in the preceding section. The "weak" interactions can be understood in terms of magnetic barrier penetration (Fig. 1), and hence in terms of *apparent* intermediaries such as $W^{\pm}$. Hadronic states are interpreted as positive-energy resonances in the associated deep magnetic wells (the 'strong' forces), and their interactions involve exchange and rearrangement of constituents (thereby explaining why secondary decay modes exist).

But electromagnetic interactions conserve parity, which is not conserved in certain weak processes. We note, however, that parity is not necessarily an intrinsic property of individual particles, but of particle *states*, and these may or may not be eigenstates of P and CP. All such parity-violating processes seem to involve neutrinos. Although a detailed theory of parity nonconservation, and a reason why we observe only left-handed neutrinos, are still elusive, we shall presume the answers to be found eventually within proposition (2) above. (This is in contrast to Fermi's and subsequent theories in which parity is violated by the *interaction* that supposedly *creates* $\bar{\nu}$.)

As noted earlier, a potential difficulty with building the unstable higher-mass particles from stable leptons is a need to explain the small magnetic moments of the former. A possible solution has been provided by Barut and Bracken,[35] in which the origin of magnetic moment is studied through a re-examination of electron *Zitterbewegung*. They find the spin magnetic moment operator for a moving electron to be

$$\boldsymbol{\mu} = ec\mathbf{S}H^{-1} \tag{4}$$

where H is the free-electron Hamiltonian. At rest $H \sim m_e$, but in a high-energy resonance state H is closer to the mass of the resonance. This energy dependence of the magnetic moment has not yet been observed directly, of course, but one might be tempted to take the observed moments of unstable particles as evidence!

3.1. Leptons

The most familiar coupling among the stable leptons is positronium, e^+e^-, in which the interaction is almost completely Coulomb with energy on the order of electron volts. At much higher energies the magnetic interactions between particles, including anomalous moments, lead to a rich new spectrum of resonances, to which we return presently.

Of crucial significance to the developing model is the pure magnetic system $\nu\bar{\nu}$, which by analogy we call *neutrinium*. This pair plays the role of a "magnetic photon," and is envisioned as being produced copiously in a $J=1$ state, with $S=0$, $l=1$. Similarly, $e^-\bar{\nu}$ and $e^+\nu$ are magnetic resonances in a $J=0$ state with $S=1, l=1$ (for -1 relative parity). All three couplings are relatively weak 2-body magnetic resonances arising in pair-production processes; with the possible exception of $\nu\bar{\nu}$, they are rather short-lived. We conjecture, however, that the 3-body and higher states are more strongly bound, much as is the case in some atoms: diatomic beryllium, for example, is only weakly bound, but binding energy per particle increases dramatically as higher-order clusters are formed.

A number of studies of these basic 2-body interactions have been carried out, going back to the work of Carlson and Oppenheimer[9] and Bethe.[10] Subsequently, detailed magnetic cross sections have been calculated for electron–neutrino scattering and neutrino pair production through e^+e^- scattering,[36,37] and for photon–neutrino processes.[38] It appears that $e^+e^- \to \nu\bar{\nu}$ is extremely difficult to detect in the laboratory. Neutral-current processes have received much attention within the electroweak model, of course, so that there is some importance to comparing the two mechanisms with one another, as well as with experiment. The angular distributions in νe scattering are quite different,[39] but the data are not yet adequate to distinguish between the two.

As we move in energy above the $e^\pm$ pair-production threshold at an MeV, the lepton spectrum emerges, owing to the binding of $\nu\bar{\nu}$ pairs to $e^\pm$, and the simplest states to be encountered are the muons: $\mu^\pm = e^+\nu\bar{\nu}$, $\mu^- = e^-\bar{\nu}\nu$. Because of the avowed fundamental role of $\nu\bar{\nu}$ as the "glue" of particle physics, the exact structure is not quite this simple. What one actually sees, for example, is a $J = \frac{1}{2}$ resonance at ~ 105 MeV with lifetime $\sim 10^{-6}$ s:

$$\begin{aligned} \mu^- &\to e^- + \bar{\nu} + \nu & &(98.6\%) \\ &\to e^- + \bar{\nu} + \nu + \gamma & &(1.4\%) \end{aligned} \tag{5}$$

What one does *not* see, though, is $\mu \to e + \gamma$, which is naturally *not* to be expected under pure magnetic binding. Thus, it is possible that the muon is actually $\mu^- = e^-(\bar{\nu}\nu)(\nu\bar{\nu})$, say, and that the preferential rearrangement is

$$\mu^- = e^-\bar{\nu}\nu_\mu, \qquad \nu_\mu \equiv \nu(\nu\bar{\nu}) \tag{6}$$

This introduces what *looks like* a new mu neutrino, ν_μ. But the two neutrinos are not observed in μ-decay, and it is certainly not known if ν_μ decays, so it is possible that the mu neutrino is an additional stable particle. Although either scenario may yet prove correct, we shall here adopt the definition of Eq. (6) as being the more economical. The ν_μ has exactly the same quantum numbers as ν, with the exception of an elusive quality called "mu-ness"—we return to a discussion of their differences presently. The suggestion is strong that there is an entire lepton spectrum to be obtained (in principle) by adding $\nu\bar{\nu}$ pairs to e and ν. For example, at ~ 1.8 GeV one finds a $J = \frac{1}{2}$ resonance $\tau^+ \equiv e^+\nu\bar{\nu}_\tau$, or

$$\begin{aligned} \tau^- &\to e^- + \bar{\nu} + \nu_\tau & &(16.2\%) \\ &\to \mu^- + \bar{\nu} + \nu & &(18.5\%) \\ &\to \text{hadrons}^- & &(65\%) \end{aligned} \tag{7}$$

with lifetimes $\sim 10^{-13}$ s. Although it has not been seen, presumably $\nu_\tau \equiv \nu(\nu\bar{\nu})(\nu\bar{\nu})$. There is so much energy available here from $e^\pm$ annihilation that almost anything can (and does) happen, as indicated in the dominant third line of Eq. (7). It is known, of course, that this is a very prolific mechanism for producing $\mu^\pm$ pairs, as well as $\tau^\pm$, and that the QED pointlike cross section $4\pi\alpha^2/3E_{\rm cm}^2$ is followed very well through ~ 40 GeV. This observation plays a significant role in later discussions at higher energies.

Barut has suggested that the lepton mass spectrum can be described by the (empirical) expression[40,41]

$$m_l = m_e \left[1 + \tfrac{3}{2}\alpha^{-1} \sum_{n=1} n^4 \right] \tag{8}$$

although one should expect relativistic corrections of the form $n^6 \to n^4/(n^2 + \alpha^2)$.
Hence,

$$\begin{aligned} m_e &= 0.511 \text{ MeV} && (\text{stable}) \\ m_\mu &= 105.55 \text{ MeV} && (\sim 10^{-6}\text{ s}) \\ m_\tau &= 1786 \text{ MeV} && (\sim 10^{-13}\text{ s}) \\ m_\delta &= 10.29 \text{ GeV} && (?) \\ &\vdots \end{aligned} \tag{9}$$

These values are in excellent agreement with experiment through m_τ; the δ is only predicted at this point. In accordance with the above discussion, note that n appears to count the number of $\nu\bar{\nu}$ pairs. Indeed, one gains the impression that "muon-ness," or muon number, simply has to do with the number of $\nu\bar{\nu}$ pairs attached to e or ν, thereby explaining the apparent lepton family structure. (Incidentally, the magnetic model also provides a natural expression for the Zel'dovich–Sakharov hadron mass formula.[42])

From an S-matrix point of view, we can adopt a model of the muon as a 3-body resonance in its decay channel, as noted in Eq. (6). The corresponding relativistic 3-body problem is quite difficult to solve, but a tractable model can be constructed by first coupling the neutrinos into a spin-1 state and then considering the Dirac equation for the electron in the presence of this field.[43] Were ν $(= \nu_e)$ and ν_μ identical, this would give a total magnetic moment zero, so that it is now essential that $\nu \neq \nu_\mu$. If the observed muon lifetime is coupled to the presumption of a ν moment $\sim 10^{-10}\,\mu_{\rm B}$, one finds a moment of similar magnitude for the ν_μ—and in both cases the spin and moment are antiparallel. In addition, the extent of the wavefunction in this model is $r \sim 10^{-23}$ cm, thereby confirming the pointlike structure of $\mu^\pm$.

3.2. Mesons

With the exception of the very special case $\nu\bar{\nu}$, the lepton series consists of $J=\frac{1}{2}$ fermions. At relatively low energies these particles can bind weakly among themselves, as in muonium (μ^+e^-)—but this is totally a Coulomb problem. Magnetic binding of type $e\nu$ is rather weak by itself, and as the energy increases to $\sim 10^2$ MeV one begins to see the importance of $\nu\bar{\nu}$ pairs. At this level a number of $J=0$ boson-like resonances begin to appear, called *mesons.* At 140 MeV the pions emerge as the spin-0 magnetic resonances

$$\pi^- = \mu^- \bar{\nu}_\mu, \qquad \pi^+ = \mu^+ \nu_\mu \tag{10a}$$

with lifetime $\sim 10^{-8}$ s; at 135 MeV there is a neutral pion with lifetime $\sim 10^{-16}$ s,

$$\begin{aligned} \pi^0 &= \frac{1}{\sqrt{2}}(e^+e^- - \nu\bar{\nu}) \\ &\to 2\gamma \qquad\qquad (\sim 98.8\%) \\ &\to e^+e^-\gamma \qquad\quad (\sim 1.2\%) \end{aligned} \tag{10b}$$

It is important here to note that the decay $\pi \to e\nu$ is almost never seen, again emphasizing the weakness of the $e\nu$ binding *in the free state.* Rather, the $\nu\bar{\nu}$ "glue" is needed to produce pions, the general structure being $\pi^- = e^-\bar{\nu}(\nu\bar{\nu})(\nu\bar{\nu})(\nu\bar{\nu})$, for example, and the preferred rearrangement is that of Eq. (10a). This is also emphasized by the extraordinarily short lifetime of π^0. Thus, far from being a curiosity with no apparent purpose, the role of the muon is essential to the structure of mesons. Ironically, the muon was first called a "mu meson" upon its discovery, for it was thought to be the meson predicted by Yukawa. Discovery of the pion eventually showed them to be fundamentally different particles, yet here we see them in a more intimate relationship.

Evidence for the structure $\pi = \mu\nu_\mu$ is quite strong—not only from the principal decay modes, but also from the way ν_μ itself was discovered. Moreover, observation of muon capture in the reaction $\mu^- + p \to n + \nu_\mu$ following formation of a muonic atom is in complete agreement with this picture of the pion, as well as with the β-decay structure of the neutron (see below). It will be useful later to note that, because the pion is a spin-0 resonance of stable leptons, the absence of a magnetic barrier permits a closer approach to nucleons than might otherwise be expected. Given the pion lifetime from experiment, a crude model in the spirit of that above for the muon has been employed to calculate the Fermi coupling constant, and to predict a magnetic moment for the neutrino comparable to the experimental upper bound.[43]

The known meson sequences are now built as $l\bar{l}$ pairs, a procedure also followed by the quark model. Mixing effects are inevitable, however, which can be seen as follows[44]: the quartet (e, ν, μ, ν_μ), viewed as a representation of $U(4)$, can be decomposed with respect to the $SU(3)$-subalgebra in two different ways, l and l', depending on whether μ or ν_μ is taken as a singlet. Therefore, the mesons can be constructed as $(l \otimes \bar{l}) \otimes (l' \otimes \bar{l}')$, a scheme that will prove quite useful when discussing internal symmetries presently.

This construction accounts for the great many short-lived resonances through 2 GeV. Most important is the fact all masses and lifetimes can in principle be calculated directly owing to the known dynamical mechanisms involved, though the many-body problems to be solved are indeed formidable. Note that all the decays involving leptons are "weak," in a sense to be compared presently to other hadronic decays. At higher energies there are a few surprises, particularly from e^+e^- scattering—and some processes appear that prevent us from thinking that mesons are nothing but inflated leptons.

One surprise that appears early in the meson spectrum is the K-meson system, which at first glance is just a high-energy version of the pion family. At ~ 494 MeV and $J=0$ we find

$$K^+ = \mu^+ \nu_\mu, \qquad K^- = \mu^- \bar{\nu}_\mu \tag{11}$$

with lifetimes $\sim 10^{-8}$ s, and a secondary decay branch $K^\pm \to \pi^\pm \pi^0$ of $\sim 21\%$. The neutral kaon appears at ~ 498 MeV and $J=0$, giving muonium-like resonances

$$K^0 = e^- \mu^+, \qquad \bar{K}^0 = e^+ \mu^- \tag{12}$$

But K^0 production in association with other hadrons indicates that there is not a direct conjugate relation between K^0 and $\bar{K}^0$. (There is also a difference of 2 in their strangeness quantum numbers, to be discussed below.) Rather, the $K^0\bar{K}^0$ system is composed of a short-lived component K_S^0 decaying to two pions in $\sim 10^{-10}$ s, and a long-lived component K_L^0 decaying principally to an odd number of pions in $\sim 5 \times 10^{-8}$ s. We note for later reference that K_S^0 never decays to a muon, and that the decays $K_L^0 \to \mu^+\mu^-$, $K^+ \to \pi^+ \nu\bar{\nu}$ are suppressed.

As is well known,[45] an understanding of this system lies with the observation that the structure of K^0 and $\bar{K}^0$ differs markedly from the decay products of K_S^0 and K_L^0. We thus *define* the latter two as eigenstates of CP,

$$K_S \equiv K^0 + \bar{K}^0, \qquad K_L \equiv K^0 - \bar{K}^0, \tag{13}$$

up to normalization. Since K_S is CP-even it decays only to two pions, and since K_L is CP-odd it can decay only to $\pi^+\pi^-\pi^0$ (say). Lifetime and mass differences now follow readily. But a small 2π-decay of K_L was discovered in 1964, so that CP is *not* conserved in this system. Both the smallness and the rarity of this violation have been a puzzle for years. The standard model is able to accommodate CP violation via complex couplings in the weak charged current, and predicts a similar effect in the B-meson system.[46]

A transparent and rather satisfying description of this system can be developed within the present particle picture.[31] Return for a moment to the "bare" model suggested by Eq. (5). Then suppose that the state $e^-\mu^+$ is produced, and envision the $\nu\bar{\nu}$ pair as being exchanged between e^+ and e^-, so that the neutrino pair oscillates back and forth in a magnetic double potential well. Such a scenario will produce K^0 and $\bar{K}^0$ alternately, and Eq. (13) provides the requisite symmetric and antisymmetric combinations. This description is formally identical to that for the NH_3 molecule, of course, and one can extract results immediately from that formalism. As examples, the differences in masses and lifetimes are reproduced quite well. Although the combinations represented by K_S and K_L are eigenstates of CP, there can arise a small violation of these symmetric and antisymmetric combinations. Re-introduce the distinction between ν and ν_μ, as in Eq. (6), so that explicitly

$$K^0 = e^-\bar{\nu}_\mu \nu e^+, \qquad \bar{K}^0 = e^-\bar{\nu}\nu_\mu e^+ \tag{14}$$

If there is now a further interaction converting $\bar{\nu}\nu_\mu$ into $\bar{\nu}_\mu\nu$, and vice versa, then a further asymmetry arises between K_L and K_S. Exchange of $\nu\bar{\nu}$ pairs between neutrinos provides an obvious mechanism. But then CP violations are predicted immediately in other neutral-meson systems of the forms $(e^-\tau^+, e^+\tau^-)$ and $(\mu^-\tau^+, \mu^+\tau^-)$, and perhaps the puzzle disappears.

3.3. Baryons

The existence and dominant role of the proton in the structure of matter as we know it would seem to present an obstacle to a pure lepton theory. But, while p has always appeared to be absolutely stable, deep inelastic scattering experiments provide incontrovertible evidence that the proton possesses an internal structure, readily ascribed to three internal constituents. Eventually, these "partons" were associated with the quark model and the three "flavors" (u, s, d). The binding is extraordinarily tight and stable at just under a GeV, with lifetime at least $>10^{30}$ years. Hence, for present purposes we can take p as absolutely stable—the only such baryon (or hadron), along with $\bar{p}$.

Feynman recalls the following fragment of a conversation with John Wheeler[47]:

> "But, Professor," I said, "there aren't as many positrons as electrons."
> "Well, maybe they are hiding in the protons, or something," he said.

A similar suggestion had been made by Oliver Lodge much earlier.[48] Thus, it is tempting to postulate that $p \equiv e^{+}(e^{+}e^{-})$, so that we have a pure lepton theory after all! A superstable, 3-body magnetic bound state of this kind has always presented the problem of too big a magnetic moment. Whether or not this is true is at least reduced to calculation here, and we believe that the mechanism of Eq. (4) might lead to the appropriate nuclear magneton. This speculation is rather liberating, for it frees us from worrying about where the antimatter is hidden—the universe can indeed be considered symmetric in that respect. This still begs the question of initial conditions, of course, for the primordial binding could as well have been $e^{-}(e^{-}e^{+})$. Nothing that follows depends on this model for the proton, as long as it is considered completely stable.

The first (in energy) baryon to be built from p is the neutron, which we interpret as $n = pe^{-}\bar{v}$ in the decay channel of its constituents. That is, we take precisely the opposite view from Fermi. Although $e^{-}\bar{v}$ is only weakly bound magnetically, and the resonance is not observed independently in the free state, the 3-body state $pe^{-}\bar{v}$ must be bound reasonably tightly—at least well enough to last ~ 15 min in the free state. Beta decay is now seen to be just penetration through a magnetic barrier of the type indicated in Fig. 1, and is a model for all other "weak" decays. While the antineutrino helps compensate for the electron magnetic moment, we still rely on the mechanism of Eq. (4) to reduce the moment of the resonance to the observed value. Note that at high energies (~ 90 GeV) it is quite possible to form a very-short-lived resonance $W = ev$ in a $J = 1$ state. But it does not seem essential, as the barrier penetration itself should be of sufficiently short range.

This model of the neutron also provides a direct explanation of the mass difference $\Delta m = m_p - m_n$, which experimentally is ~ -1.29 MeV. It had been long thought that Δm should be positive and electromagnetic in origin, but failure to obtain the correct sign has tended to move the search for an explanation to higher energies. If n is a positive-energy magnetic resonance, however, there is no surprise that $\Delta m < 0$, and all that remains is to carry out the 3-body calculation to verify the magnitude.

The preceding discussion suggests an interpretation of nuclear interactions as an exchange of ev pairs between protons. Similar suggestions had been made long ago by Tamm[49] and Iwanenko,[50] but using weak interactions and the Fermi coupling, rather than magnetic forces. Barut has proposed a model of the nucleus as a close-packed lattice of A protons with N quasi-free $e\nu$ pairs,[31] which tends to reconcile the contradictory features of the shell and liquid-drop models. It also introduces the meson theory of nuclear forces as an immediate approximation, a possibility not contained even in principle in the quark model. Although $e^-\bar{\nu}$ is not quite the same as a free π^-, it does form its skeletal structure and possesses the same basic quantum numbers. Because it is a spin-0 resonance, the pion can penetrate magnetic forces deeply and so provide nuclear binding. This is the "strong" interaction. Despite the weakness of the $e\nu$-binding, it is nevertheless in line with nuclear binding energies, which are never more than 10 MeV per particle.

We see, then, that the basic baryon is $p(l\bar{l})$—the general baryon is p with a cloud of $l\bar{l}$ pairs, and ultimately decays to these constituents.

3.4. Internal Symmetries and Hadron Multiplets

In this already-unified particle picture one attains complete understanding of the so-called internal symmetries and associated quantum numbers, though a brief tour of the conventional view will be useful in setting the context. The apparent charge independence of nuclear forces led to introduction of the notion of *isospin*, I, and in analogy with the angular-momentum formalism the third component I_3 takes $(2I+1)$ values. Thus, n and p are just two manifestations of a single entity, the nucleon, which has $I=\frac{1}{2}$. Electromagnetic interactions break this symmetry into the two states $I=\pm\frac{1}{2}$. The nucleon is therefore an $SU(2)$ doublet with charge $Q=I_3+\frac{1}{2}$ (in units of e). Similarly, the pion system with $I=1$ is an $SU(2)$ triplet with $Q=I_3$, and hadrons are now described by the two internal quantum numbers (Q, I).

As the energy scale moved above a GeV, and up through 2.5 GeV, many new baryon resonances appeared through the "strong" interactions typified by the short-range nuclear force. While most of these decayed strongly, with lifetimes $\sim 10^{-23}$ s, a number of particles exhibited strange behavior by decaying much more slowly, on the order of weak-interaction lifetimes ($\sim 10^{-10}$ s). This series of "strange" particles (Λ, $\Sigma^{\pm}$, Σ^0, Ξ^0, Ξ^-, Ω^-) also exhibited primarily nonleptonic decays. The observation that these particles were always created in pairs (associated production) led to a convention for assigning a new *strangeness* quantum number S, which is conserved in strong interactions. When S was coupled with the notion of baryon number ($B=\pm 1$ for baryons and antibaryons, respectively, and is 0 for mesons and leptons), the Gell-Mann–Nishijima formula emerged,

$$Q=I_3+\tfrac{1}{2}(B+S) \tag{15}$$

leading eventually to the $SU(3)$-multiplet structure and the quark model. (One conventionally now refers to just *hypercharge*, $Y = B + S$.) But no deeper insights into the physical significance of these quantum numbers was forthcoming.

Let us now examine this organizational scheme in the context of the general picture being advocated here, in which we ought to be able to understand things from a dynamical basis—namely, pair processes and magnetic interactions among stable particles. Thus, lepton number is automatically conserved and, on the time scales considered here, the proton is stable and baryon number is always conserved. Charge is also conserved, by definition of the model. In addition, $\mu^{\pm}$ play an important role in the structure of hadrons and are stable on hadronic time scales, so we shall include them in the definitions of the quantum numbers. If n_a is the number of particles of type a present in some process, it is convenient to define the difference $N_a \equiv n_a - n_{\bar{a}}$. Then, aside from the energy, the only truly conserved quantum numbers in this picture are

$$\begin{aligned} Q &\equiv N_p + N_e + N_\mu \\ B &\equiv N_p \\ L &\equiv N_e + N_\nu + N_\mu \end{aligned} \tag{16}$$

The meson quantum numbers follow from those of $l\bar{l}$ states, as discussed above, giving vector as well as pseudo-scalar mesons. Baryon states are constructed as $p \otimes l \otimes \bar{l}'$. On hadronic time scales the muon behaves like an electron, forming magnetic pairings and resonances. The neutron is not quite a μ bound to p—having less constituents perhaps renders it more stable. But all other hadrons contain muons, so we interpret associated production as $\mu^+\mu^-$ pair production, along with copious neutrino pairs. (On this level the muons are considered stable.) For $S \neq 0$, strangeness simply counts the number of $\mu^{\pm}$ pairs in hadrons:

$$S \equiv N_\mu \tag{17}$$

Pions are not produced this way—in fact, pion–nucleon collisions are used to produce strange hadrons—so a value $S = 0$ is assigned to pions and nucleons.

Clearly, strangeness is conserved in strong production, as well as in strong decays (which are essentially restricted to very short-lived baryon resonances). These latter decay so fast that there is little time for any rearrangement to take place. But in the longer-lived strange hadrons muon decay is partially suppressed relative to the free state, so that among the hyperons we see mostly semileptonic decays to other hadrons—usually one

$\mu^{\pm}$ decays in a strangeness-changing decay of $|\Delta S| = 1$. Note that this automatically accounts for the selection rule $\Delta Q = \Delta S$. It now seems that the Cabibbo angle simply measures the suppression factor for muon decay inside hadrons. In a similar manner, we can identify *charm* as $C \equiv N_{\nu_\mu}$, which presumably explains the suppression of $K^+ \to \pi^+ \nu\bar{\nu}$, for example. Note carefully that all these internal quantum numbers are unaffected by addition of e^+e^- and $\nu\bar{\nu}$ pairs to any state.

Interpretation of the nuclear interaction as exchange of $e^-\bar{\nu}$ between two protons leads immediately to an isospin description based on quantum-mechanical exchange symmetry alone. (This is analogous to e^--exchange in H_2^+.) For two nucleons the states of definite isospin are

$$\begin{aligned} I = 1&: |pp\rangle, |nn\rangle, \frac{1}{\sqrt{2}}(|pn\rangle + |np\rangle) \\ I = 0&: \frac{1}{\sqrt{2}}(|pn\rangle - |np\rangle) \end{aligned} \tag{18}$$

For two pions, similar exchanges yield states $|\pi^{\pm}\pi^{\pm}\rangle$, $\frac{1}{\sqrt{2}}[|\pi^{\pm}\pi^0\rangle + |\pi^0\pi^{\pm}\rangle]$, etc., and from the Pauli principle one now understands the spin dependence of strong interactions.

Isospin symmetry is nothing more than symmetric or antisymmetric exchange or rearrangement of stable constituents among hadrons, and the third component just counts the number of stable constituents:

$$I_3 \equiv \tfrac{1}{2}(N_p + N_e + N_\nu) \tag{19}$$

From Eqs. (16), (17), and (19) one verifies the Gell–Mann–Nishijima formula (15)—for charmed particles C is added to the right-hand side when the level is such that ν_μ is approximately stable.

Conservation of isospin and strangeness are only approximate symmetries, of course, arising from the resonance structure of hadrons in terms of stable particles. Traditionally, they have been perceived to have deeper significance and incorporated into continuous group structures, such as $SU(3)$, and a multiplet structure is revealed by plotting S (or Y) versus I_3. But, as already mentioned, the same group structure is readily obtained from the magnetic model of stable particles.[32,44,51] As examples, the pseudo-scalar meson octet ($J^P = 0^-$) and the baryon octet ($J^P = \frac{1}{2}^+$) and decouplet ($J = \frac{3}{2}^+$) are illustrated in Fig. 2. It is clear that additional $\nu\bar{\nu}$ pairs are acquired, either in production or decay, to yield observed products. Note again, however, that arbitrary numbers of $\nu\bar{\nu}$ and e^+e^- pairs can be added to each hadron without changing the internal quantum numbers. As has been known for a long time,[52] only *finite* subgroups of

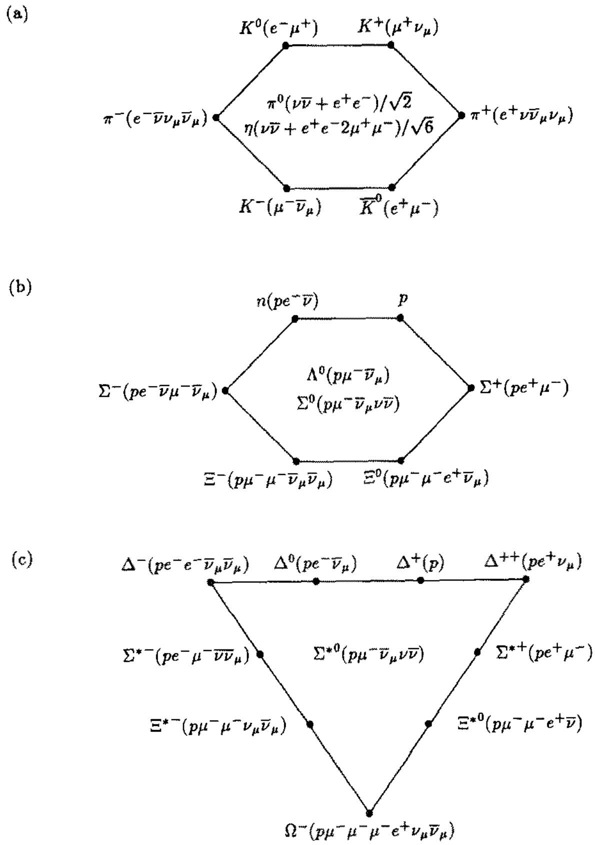

Fig. 2. Multiplet structure in the magnetic-potential model: (a) pseudo-scalar meson octet; (b) baryon octet; (c) baryon decouplet.

$SU(2)$ and $SU(3)$ are required for what now appears to be just a convenient classification scheme. Because the full, continuous Lie groups are not needed, there would seem to be no compelling reason to gauge these groups. In connection with the Yang–Mills non-Abelian gauge theory, "...conservation of isotopic spin only suggests, and does not require, the existence of an isotopic spin gauge field."[53]

4. SOME FINAL COMMENTS

What we have reviewed here would seem to be, at the very least, the beginnings of a desirable description of how the world works; but at best it is a program, rather than a full-blown theory. It does not provide complete unification, in that it does not include the gravitational field—but neither does the standard model. In any event, there is no evidence that gravity plays any role whatsoever at the particle level. To proceed further it is necessary to learn how to carry out detailed calculations involving few-body interactions. At present, only some two-body problems have been studied, and there is evidence that these are simply inadequate for developing the full scope of the model. This situation is exacerbated by the marginal values being reported as upper bounds for the neutrino magnetic moment, although the resonances necessary to produce the particle spectrum may be strongly dependent on 3-body forces. Additional mathematical difficulties arise from the observation that it is necessary here to consider localized wavefunctions, and to avoid the use of perturbation theory. Thus, further development of nonperturbative QED is a basic requirement for future progress. This may not be the way it works, but until all the loopholes have been closed it is difficult to ignore, and surely worthy of continued effort.

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For thirty years I have benefitted from exposure to Asim Barut's fertile imagination, and the example he sets of one who does his own thinking. It is a pleasure to dedicate this article to a celebration of his sixty-fifth birthday.

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