an upper limit on m_φ:

$$m_\varphi < 4.6\times10^{-27} m_e. \tag{9}$$

So far as elementary-particle physics is concerned, the Higgs particle is massless, a property not inconsistent with experiment.[3] (It is interesting to note that if m_φ arises gravitationally, then $\lambda \sim G \sim \mu^2$ and the necessary factor of smallness arises here but in a more intricate fashion than previously surmised.[4])

There are a number of consequences which ensue. If it be granted that the electron mass originates from spontaneous symmetry breakdown via the Higgs mechanism, then φ is weakly coupled to the electron with coupling constant

$$G_\varphi \sim 2\times10^{-6}. \tag{10}$$

This small coupling inhibits easy production and makes it plausible that the particle has escaped detection. If the Higgs particle is effectively massless, however, it must also couple to the proton with a sign opposite to the electron coupling, otherwise matter would collapse under this weak but coherent long-ranged "scalar" electrodynamic force. One place to look for the particle is in the $0^+ \to 0^+$ transition in ^{16}O.

If the assumption is made that the universe at the present epoch is isotropic, a second consequence of the model, somewhat more speculative, ensues. The appearance of a term with the cosmological constant in the equations governing the evolution of the universe implies that the universe will eventually contract.[6] This follows from the property that the cosmological constant is negative [Eq. (8)] independent of the φ mass as long as it is nonzero. A nonzero value of m_φ is required for the Higgs mechanism to work in the usually assumed manner. It seems striking that the absence of both nonrenormalizable ultraviolet divergences in weak interactions and a divergent expansion of the universe might have something in common. If T nonconservation originates from spontaneous breakdown of symmetry as, for example, in the Lee model,[7] then the Higgs field could conceivably also put the nonconservation in a cosmological context.

The discovery of a nearly massless scalar field with a weak coupling is a necessary condition for the validity of the hypothesis that the Higgs field is also a cosmological field.

I hope to report more detailed consequences of the cosmological hypothesis in the near future.

I wish to thank my colleague, Professor K. Mahanthappa, for his advice and criticisms.

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Scalar Models of Weak Interactions*

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A model of weak interactions mediated by scalar bosons $B^\pm$, B^0, and $\bar{B}^0$ is presented.

The recent discovery[1-4] of muonless neutrino scattering events has lent strong credence to the essential correctness of some form of the Weinberg-Salam theory[5,6] of weak interactions. It seems particularly important to see now whether the data can be fitted by other theories as well; in this spirit I wish to discuss a class of renormalizable models of weak interactions. They are basically elaborations of models discussed years ago,[7-11] modified by a few new observations.

The characteristic feature is that the weak interactions are mediated by scalar bosons, $B^\pm$, B^0, and

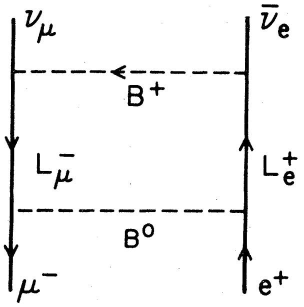

FIG. 1. Diagram for μ decay.

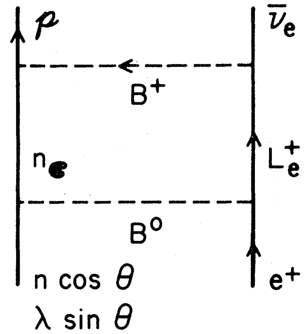

FIG. 2. Diagram for β decay.

$\bar B^0$ (for the moment assume B^0 is not self-conjugate). In addition to the usual leptons e^-, ν_e, and μ^-, ν_μ I also introduce two massive charged leptons L_e^- and L_μ^-; the baryons are assumed to be made up of quarks, $\mathcal{P}$, $\mathcal{N}$, λ, and the charmed quark $\mathcal{P}'$,[12] where the notation of $\mathcal{N}_c = \mathcal{N}\cos\theta + \lambda\sin\theta$ and $\lambda_c = -\mathcal{N}\sin\theta + \lambda\cos\theta$ will also be used, θ being of course the Cabibbo angle. I claim that no additional particles are necessary.

The interaction Lagrangian is

$$\mathcal{L}_{\rm int} = -if\Big\{\sum_{l=e,\mu}\bar L_l(1-\gamma_5)lB^0 + \bar L_l(1-\gamma_5)\nu_l B^- + [\bar{\mathcal{N}}_c(1-\gamma_5)\mathcal{P} + \bar\lambda_c(1-\gamma_5)\mathcal{P}']B^- + [\bar{\mathcal{N}}_c(1-\gamma_5)\mathcal{N}_c + \bar\lambda_c(1-\gamma_5)\lambda_c]B^0\Big\} + \text{H.c.} \tag{1}$$

The diagram for μ decay is shown in Fig. 1. As is well known, this leads to an effective $V-A$ interaction {remember $[\gamma_4(1-\gamma_5)]^\dagger = \gamma_4(1+\gamma_5)$}

$$H_{\rm eff} = M_B^{-2}(f^2/4\pi)^2\bar\nu_\mu\gamma^\alpha(1-\gamma_5)\mu\bar e\gamma_\alpha(1-\gamma_5)\nu_e \tag{2}$$

plus corrections, which are $O(M_\mu^2/M_B^2)$. In evaluating the diagram of Fig. 1, I have assumed M_B to be much larger than the mass of L_e^- or L_μ^-, which in turn is taken to be greater than or equal to ~5 GeV. This last limit is imposed by the present limit[13] on the production of charged massive leptons which decay into neutrino plus hadrons. In this model $\bar\nu_\mu + e^- \to \bar\nu_\mu + e^-$ is forbidden.

The conventional value of μ decay follows if we take

$$\left(\frac{f^2}{4\pi}\right)^2\frac{1}{M_B^2} = \frac{G_F}{\sqrt2} = \frac{1}{\sqrt2}\,\frac{10^{-5}}{M_{\rm proton}^2}. \tag{3}$$

For definiteness let us take $M_B \sim 300$ GeV so that $f^2/4\pi \sim 1$.

The usual semileptonic weak interactions with the Cabibbo angle are as depicted in the diagram of Fig. 2. Of course, with $f^2/4\pi$ being of order 1, universality might be destroyed by higher-order terms. I have verified that this does not occur in the one-loop correction to vertices if the Lagrangian (1), is used. A typical one-loop vertex correction is displayed in Fig. 3.

Turning to $\Delta S = 1$, $\Delta Q = 0$ semileptonic decays, we find the model compatible with present limits. Since the hadronic coupling to B^0 may be rewritten as

$$-if\{\bar{\mathcal{N}}(1-\gamma_5)\mathcal{N} + \bar\lambda(1-\gamma_5)\lambda\}B^0, \tag{4}$$

we see that $\bar{\mathcal{N}}\lambda \not\to \mu^+\mu^-$ by two-B^0 exchange and hence $K_L^0 \not\to \mu^+\mu^-$, and of course $K^+ \not\to \pi^+e^+e^-$ to order G_F. $K^+ \to \pi^+\bar\nu\nu$ or rather $\bar\lambda\mathcal{N} \to \bar\nu\nu$ is allowed by B^+, B^- exchange as shown in Fig. 4, but the well-known cancelation between $\mathcal{P}$ and $\mathcal{P}'$ leads to an additional factor, $\Delta m^2/M_B^2 \sim G_F$, if we take the square of the $\mathcal{P}$ - $\mathcal{P}'$ mass difference to be of the order of a few GeV^2 and $M_B \sim 300$ GeV.

Turning now to nonleptonic weak interactions we find that the diagrams of Figs. 5(a) and 5(b)

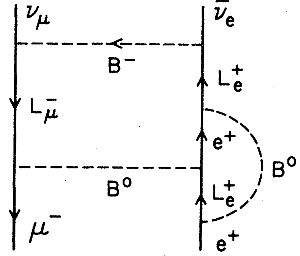

FIG. 3. Typical one-loop weak correction to μ decay.

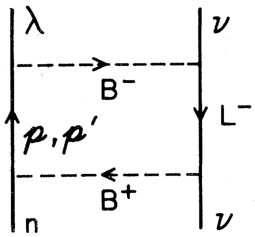

FIG. 4. Diagram for $\mathcal{N}\bar{\lambda}\to\nu\bar{\nu}$.

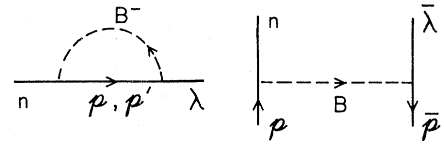

FIG. 5. Nonleptonic $\Delta S=1$ transition.

lead to nonleptonic $\Delta S=1$ amplitudes roughly a factor of 10 larger than semileptonic $\Delta S=1$ allowed transitions. This is, I believe, not incompatible with experiment, though a more detailed dynamical calculation might lead to difficulties. $\Delta S=2$ transitions are suppressed of course by the $\mathcal{P}$-$\mathcal{P}'$ charm cancelation mechanism. A typical lowest-order allowed diagram for the $\Delta S=2$ transition is shown in Fig. 6 and is of order $(\Delta m^2/M_B{}^2)(f^2/4\pi)^2\sim G_F{}^2$ because of the two $\mathcal{N}\to\lambda$ transitions which are necessary if we wish to have $\Delta S=2$, $\Delta Q=0$.

The $\Delta S=0$ hadronic weak interactions appear to present considerable difficulties. The diagrams of Figs. 7(a) and 7(b) are in fact logarithmically divergent because there is no cancelation mechanism so that on the surface one has large violations of isospin and parity. The diagrams involving B^0 and $\bar{B}^0$ have canceling parity-nonconserving parts. To the original Lagrangian one could add a term

$$\mathcal{L}_{\rm int}\to\mathcal{L}_{\rm int}\pm if\{\bar{\mathcal{P}}(1-\gamma_5)\mathcal{P}+\bar{\mathcal{P}}'(1-\gamma_5)\mathcal{P}'\}B^0+\text{H.c.} \qquad (5)$$

This gives a contribution to β decay which is $O(m_L{}^2/M_B{}^2)$ because of $1\pm\gamma_5$ factors and also does not modify the estimate of one-loop corrections to β decay. It does make the B^0, $\bar{B}^0$ loops as shown in Fig. 7(b) isospin invariant as well as parity invariant. But what about the $B^\pm$ loops? The leading logarithmically divergent behavior is an effective wave-function renormalization and can be transformed away, as explained by Low.[14] The finite part is proportional to (k is a quark momentum)

$$\frac{f^2}{4\pi}\frac{k^2}{M_B{}^2}\not{k}(1\pm\gamma_5), \qquad (6)$$

but $f^2/4\pi M_B{}^2\sim G_F$ and so we need not worry about (6).

Neutral $\Delta S=0$ hadronic currents are present in this model; the lowest-order diagrams for ν_μ $(\bar{\nu}_\mu)$ + hadron $\to\nu_\mu$ $(\bar{\nu}_\mu)$ + hadron are depicted in Figs. 8(a) and 8(b). The resultant effective neutral-current Hamiltonian, using (3), is

$$H_{\rm eff}=(G_F/\sqrt{2})\bar{\nu}\gamma^\alpha(1-\gamma_5)\nu\{\bar{\mathcal{N}}\gamma_\alpha(1+\gamma_5)\mathcal{N}+\bar{\lambda}\gamma_\alpha(1+\gamma_5)\lambda-\bar{\mathcal{P}}\gamma_\alpha(1-\gamma_5)\mathcal{P}\}, \qquad (7)$$

so that we have an effective $\Delta I=1$ vector current and $\Delta I=0$ axial-vector current. In that sense it is different from the proposed models for neutral currents.[5,6,15,16] The predicted cross section for muonless events is presumably too large to agree with experiment by a factor of the order of 3. Agreement with experiment may be achieved (i) by multiplying all couplings to $B^\pm$ by a factor $\epsilon<1$ or (ii) by increasing the mass of $B^\pm$ relative to the B^0, $\bar{B}^0$ mass. The first scheme is similar to that of Adler and Tuan.[16] Agreement with experiment follows of course because $\nu_\mu+N\to\nu_\mu+N$ proceeds by exchange of B^+ and B^- whereas β decay and μ decay involve exchange of B^+ and B^0.

A second class of models along the same lines can be constructed by replacing the charged leptons $L_e{}^-$ and $L_\mu{}^-$ by neutral ones $L_e{}^0$, $L_\mu{}^0$. The leptonic coupling then becomes

$$\mathcal{L}_{\rm int}{}^{\rm lept}=-if\{\sum_{l=e,\mu}\bar{L}_l(1-\gamma_5)\nu_l B^0+\bar{L}_l(1-\gamma_5)lB^-\}. \qquad (8)$$

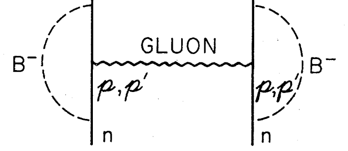

FIG. 6. $\Delta S=2$ transition.

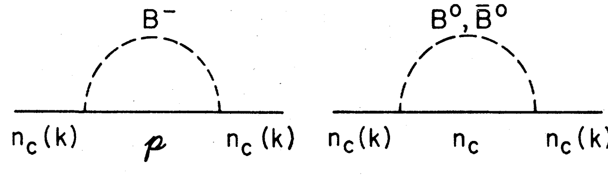

FIG. 7. $\Delta S=0$ transitions.

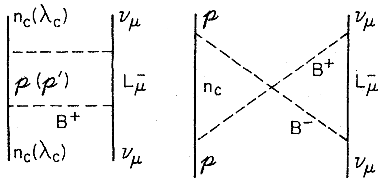

FIG. 8. Effective–weak–neutral-current diagrams for model with L_μ^-.

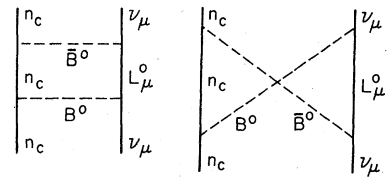

FIG. 9. Effective–weak–neutral-current diagrams for model with L_μ^0.

The main difference now is that if B^0 is self-conjugate, there is a cancelation between the direct and crossed graphs in $\nu_\mu + n_c \to \nu_\mu + n_c$ as shown in Figs. 9(a) and 9(b). If $B^0 \neq \bar{B}^0$, the arguments proceed as before except that the factor ϵ should be in the B^0 couplings and/or $M_{B^0} > M_{B^+}$.

Finally, there is one more possibility worth discussing in this context. There is basically no experimental limit on neutral heavy-lepton masses[13] so that we should consider the possibility of producing an L_μ^0. It would decay with comparable rates by two modes (more into μ^- because of the factor ϵ or M_{B^+}/M_{B^0}):

$$L_\mu^0 \to \begin{cases} \nu_\mu + \text{hadrons (charge 0)}, \\ \mu^- + \text{hadrons (charge +1)}. \end{cases} \tag{9}$$

The rapidity distributions would of course be different from the usual ones because the amplitude is $S-P$ rather than $V-A$. If the L_μ mass were low enough, one would expect this to be the dominant mode in neutrino scattering:

$$\begin{aligned} \nu_\mu + N \to\ & L_\mu^0 + \text{hadrons} \\ & \hookrightarrow \mu^- + \text{hadrons} \\ & \hookrightarrow \nu_\mu + \text{hadrons}. \end{aligned} \tag{10}$$

The amplitude would be of order f^2/M_B^2 instead of $(f^2/4\pi)^2 M_B^{-2}$ and hence would supposedly dominate. This need not continue to be true however as we let M_B become larger, so much so that $f^2 \gg 4\pi$, i.e., a strong-coupling limit. There may also be kinematical suppressions of $S-P$ with respect to $V-A$ at high energies, though this does not seem to be true for inclusive processes.[17] Letting $f^2 \gg 4\pi$ would also suppress the importance of the diagram in Fig. 5(b).

In a lengthier publication, I will present the details of these calculations and further considerations. I have still not considered the problem of going to higher orders in $f^2/4\pi$ and perhaps perturbation theory is not an appropriate tool. I have checked what I could; a breakdown of universality seems not unlikely.[18] It has been shown that the diagrams of Fig. 3 do not cause it; we have also seen, to order $f^2/4\pi$, that the overall renormalization required by the diagrams of Figs. 7(a) and 7(b) and their lepton analogs do leave universality unchanged. I would expect, however, an effect due to gluons coupling hadrons inside the box diagram to those outside the box diagram.

Theories of this class nevertheless have enough interesting features to warrant attention, even though they lack the elegance of gauge models. Renormalizability is straightforward, the Higgs mechanism and its accompanying scalar mesons become superfluous, and finally, the idea of weak interactions being "strong" from the beginning is intriguing.

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COMMENTS

Electronic Recombination of He_2^+

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The measured molecular spectrum of a helium afterglow plasma does not fit with available recombination theories. Highly excited states are in thermal equilibrium at a substantially different temperature than that of free electrons.

Recently published results from this laboratory[1] have shown that the electron density evolution under closely controlled electron temperatures between 300 and 500°K, over a limited range of neutral and electron densities, was in good agreement with the predictions of collisional-radiative recombination theory. In order to complement those results, we have undertaken a systematic spectroscopic study of the helium afterglow under conditions similar to those of Ref. 1. Although quite coherent, preliminary results do not fit well with any existing recombination theory.

Essentially, excited electronic states of the neutral molecule He_2 with principal quantum numbers $n > 12$ are in thermal equilibrium at a distribution temperature about 600°K under afterglow conditions where the free-electron temperature deduced from electron collision frequency measurements is closely equal to the gas temperature, i.e., 300°K. The neutral gas pressure was 20 Torr; the afterglow spectrum was dominated by He_2 bands, and no impurities were detected. 250 μsec after termination of the discharge pulse, the distribution temperature for high-n levels reaches an almost constant value while electron density decays from $\sim 5\times10^{11}$ to 10^{11} cm^{-3}. Gas heating resulting from the short-duration discharge pulse is negligible.

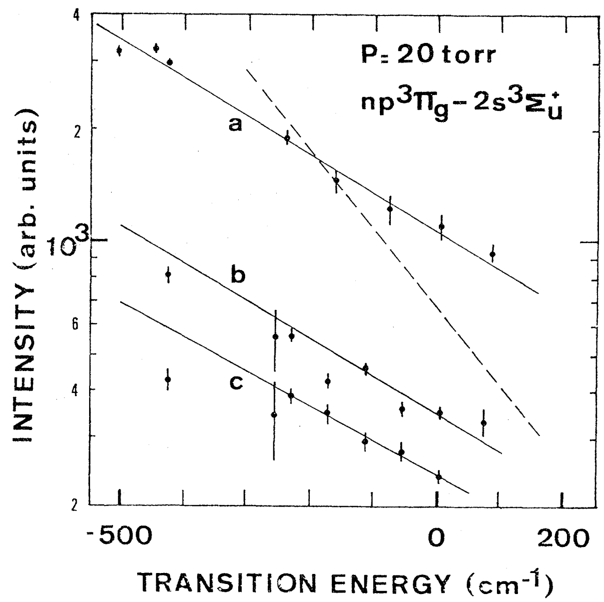

FIG. 1. Spectral intensities near the $np\,^3\Pi$-$2s\,^3\Sigma$ series limit per unit energy interval plotted versus transition energy; the origin corresponds to $He_2(2s\,^3\Sigma,\ J=0,\ v=0)$ ionization energy. Data were taken at two times in the afterglow; error bars correspond to ±1 standard deviation. a, 500 μsec, $n_e = 2.4\times10^{11}$ cm^{-3}; b, 1200 μsec, $n_e = 1.0\times10^{11}$ cm^{-3}, without microwave heating; c, electrons selectively heated at 1200 μsec, $\Delta T_e = 190$°K. The slopes yield a distribution temperature of 610°K in both a and b and 680°K in c; dashed line shows the slope corresponding to 300°K.