Experimental Restrictions on Ne'eman's Fifth Interaction

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(Received 4 August 1964)

Recently, Ne'eman has proposed a "fifth interaction" between the strangeness current and a neutral vector meson \( \chi \), for the purpose of breaking SU(3) symmetry. We show that a \( \chi \) mass less than 2\( m_{\pi} \), would be inconsistent with a variety of experiments, including \( K \)-mesonic atoms, the long-range \( pp \) potential, \( K \) regeneration from a \( K \) beam, the Lamb shift, modern refinements of the Cavendish "ice-bucket" experiment, and the absence of \( e^+ \rightarrow \gamma + \chi \) and \( \chi \rightarrow e^+ + e^- \). The remaining possibility, that \( m_{\chi} \) exceeds 2\( m_{\pi} \), is discussed briefly.

I. INTRODUCTION

RECENTLY, Ne'eman has proposed an interaction between the strangeness current and a neutral vector meson \( \chi \), which would break SU(3) symmetry in much the same way that the interaction between the electric current and the photon breaks the symmetry of isotopic-spin space. To account for the rather large violations of SU(3) symmetry, the \( \chi \) coupling to the strangeness current must have \( g_{\chi}^2 \hbar c = \gamma \) to \( \gamma \), thus providing a "fifth interaction" with strength intermediate between the strong and electromagnetic interactions. Ne'eman stressed that the fifth interaction could explain not only why SU(3) symmetry is badly broken, but also why isospin and strangeness remain good quantum numbers—a question on which the alternative mechanism of spontaneous symmetry violation in a bootstrap calculation has thus far failed to shed any light.

A particle interacting only with the strangeness current would be rather hard to detect in our non-strange world, as Ne'eman pointed out, and at first sight even a massless \( \chi \) might have escaped previous notice. This suggestion that a light particle with fairly strong coupling might have gone undetected all these years stimulated a search for experiments in which the \( \chi \) should have showed up if it exists. In the present paper we record the experimental arguments we have thought of, making no pretense that they are exhaustive. For all \( \chi \) masses in the range \( 0 \leq m_{\chi} \leq 2m_{\pi} \), we find experiments which rule out the \( \chi \) by factors of order 10, and over part of this range limits on vacuum polarization and the rate of \( K \) regeneration from a \( K \) beam, for example, rule out the \( \chi \) by much larger factors. Present experiments do not rule out \( \chi \) with mass \( m_{\chi} > 2m_{\pi} \), though the only known candidate at present is the \( \phi \), as discussed by Ne'eman.

We present our experimental cases and discussion roughly in order of increasing \( \chi \) mass, progressing from \( m_{\chi} < 10^4 \) eV (Sec. II) to \( m_{\chi} < 2m_{\pi} \) (Sec. IV) to \( m_{\chi} > 2m_{\pi} \) (Sec. V). The evidence is summarized in Table I. Section III deals with the anomalous \( K \) regeneration from a \( K \) beam observed by Leipuner et al. We find that the fifth interaction could provide an explanation of anomalous regeneration were it not for the accumulated evidence against \( m_{\chi} < 2m_{\pi} \) presented in Secs. II and IV, which makes the explanation untenable.

II. EVIDENCE AGAINST \( m_{\chi} < 2m_{\pi} \)

A. \( K \)-Mesonic Atoms

If the \( \chi \) were massless or very light, a \( K^- \) meson trapped in an atom would cascade down to the low atomic levels primarily by \( \chi \) emission rather than by \( \gamma \) emission, because the \( \chi \) coupling \( g_{\chi}^2 \gamma = \gamma \) is at least 10 times stronger than electromagnetic interactions. The time required to reach low levels where nuclear capture takes place would be reduced, and 90\% or more of the expected \( \chi \) rays would be missing. Similarly, the number of Auger electrons, emitted when the \( K \) drops in level and gives off a virtual photon that is absorbed by an orbiting electron, would be reduced by 90\% due to the which is not exactly conserved \( (\phi, J \neq 0) \) when weak interactions are taken into account. It follows that \( m_{\chi} \) cannot vanish for the \( \chi \). The argument does not seem to lead to any conclusive restriction on the physical mass, however, so we shall not make use of it.


2 Note added in proof. For the same reason the fifth interaction cannot explain the \( K \rightarrow 2 \nu \) decays observed by J. H. Christenson, J. W. Cronin, V. L. Fitch, and R. Turlay, Phys. Rev. Letters 13, 138 (1964). Note that in this latter experiment the effect of Leipuner et al. is not seen.

3 We are much indebted to Dr. William Wagner for suggesting \( K \)-mesonic atoms as a possible source of information.

4 Note that the atomic levels themselves are shifted very little by the fifth interaction because the nucleus is not strange and \( \chi \) exchange between nucleus and \( K \) occurs rarely.
TABLE I. Limits on $m_X$.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>Range of $m_X$ excluded by experiment</th>
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<tbody>
<tr>
<td>$K_\text{mesonic atoms (Sec. II-A)}$</td>
<td>$m_X &lt; 10^4$ eV by experiment</td>
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<tr>
<td>Long-range $pp$ potential (II-B)</td>
<td>Anomalous moment $K_{g\ell}/2m_X &lt; 0.03g_\ell/2m_X$ if $m_X &lt; 10^4$ eV</td>
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<tr>
<td>$K_1$ regeneration from a $K_2$ beam (II-C)</td>
<td>$m_X &lt; 2\times 10^4$ eV</td>
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<td>Lamb shift (IV-A)</td>
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<td>$10^{-10}$ eV $&lt; m_X &lt; 10^{-4}$ eV</td>
</tr>
<tr>
<td>$\chi^0 \rightarrow \gamma + \chi$ (IV-B)</td>
<td>$m_X &lt; 1.4\times 10^9$ eV</td>
</tr>
<tr>
<td>$\chi \rightarrow e^+ + e^-$ (IV-C)</td>
<td>$10^6$ eV $&lt; m_X &lt; 2.8\times 10^6$ eV</td>
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</tbody>
</table>

competition from $\chi$ emission (the $\chi$ does not couple to electrons and thus would not be reabsorbed by them).

Little experimental work on $K_\text{mesonic atoms}$ has been reported, but what does exist shows no evidence of the $\chi$. Eisenberg and Kessler have searched $K^-$ capture in emulsion for Auger electrons with energies exceeding 15 keV. Their findings are quite consistent with standard electromagnetic theory; thus they are inconsistent with the massless $\chi$ by an order of magnitude. Of course, this result does not rule out the $\chi$ if it is too massive to be emitted in $K_\text{mesonic transitions}$. In another experiment, Kopelman et al. studied how long it takes the $K$ to cascade down and get recaptured by the nucleus. They obtained along time, inconsistent with emission of massless or very light $\chi$’s. Their result does not even show signs of the expected speed-up from Stark mixing, however, so perhaps one must reserve judgment on the understanding of this particular experimental situation.

### B. Long-Range $pp$ Potential

Although nucleons have strangeness zero, they do have internal currents of $K^0$, $\Sigma^-$, etc., which allow them to couple to the $\chi$ via “strangeness form factors” and anomalous magnetic moments, just as the neutron couples to photons. At low-momentum transfers, the dominant term is the anomalous magnetic moment, and $\chi$ exchange between nucleons has the same spin dependence as the usual magnetic-magnetic interaction. If the $\chi$ is massless, the ratio of $\chi$ exchange to magnetic-magnetic photon exchange in the potential between two protons is

$$R = \left[ g^2_K K_\chi / \epsilon (1 + K_\gamma) \right]^2,$$

where $g^2_K K_\chi / \epsilon (1 + K_\gamma) = g^2_K K_\chi / 2.79 \epsilon$ is the ratio of strange anomalous moment to electromagnetic total moment. Putting in numbers, one finds $R = 1.8 K^2_\chi$.

Now the magnetic-magnetic potential between two protons at molecular distances has been checked to one part in a thousand by Ramsey, in a microwave study. From Eq. (1) we deduce that $K_\chi \lesssim 0.03$, unless $\chi$ is so massive that the factor $\epsilon^{-m_X}$ in the $\chi$-exchange potential cuts off before molecular distances are reached (i.e., $m_X \gtrsim 10^9$ eV). Of course, the limit obtained on $K_\chi$ is not an improbably low one, especially when we remember that the scalar electromagnetic moment is only $0.06 e/2m$, but it does seem rather small and will be useful later in limiting anomalous regeneration of $K_\beta$ via $\chi$ exchange between $K_2$ and nuclei.

### C. $K_1$ Regeneration from a $K_2$ Beam

Consider a $K_2$ beam approaching a bubble chamber. If $m_\chi$ is sufficiently small, the $K_2$’s will undergo long-range interactions with nuclei in the chamber by means of $\chi$ exchange. We shall study the interaction between a single nucleus and $K_2$, and then add up the coherent interactions between many nuclei and $K_2$.

In strong interactions, $K_2$ and $K_1$ are distinguished by their opposite behavior under charge conjugation, and $\chi$ is negative under charge conjugation just like the photon. Therefore, $K_2$ can turn into $K_1$ by emitting a $\chi$. The nucleus can absorb this $\chi$ by means of its anomalous “strange” magnetic moment, as discussed in Sec. II-B, or possibly through higher moments or inelastic effects. At low-momentum transfers, such as we are interested in, these latter effects can be neglected and only the “strange” magnetic moment matters.

At this point, let us recall the Pauli theory for an electron in the Coulomb field of a hydrogen atom. The effective potential in this theory includes the “spin-orbit” term

$$V_{\text{eff}} = -(1 + 2K) (\mu_0/2) \sigma \cdot E \times v$$

$$= -[(1 + 2K) / (4m^2)] \sigma \cdot r \times v,$$

where $E = \sigma / r^2$ is the electric field, $v$ is the electron velocity, $\mu_0$ is $e/2m$, and the complete magnetic moment is $(1 + K)$ times the Dirac moment. Now the $\chi$-exchange potential between $K^0$ and a nucleus is completely analogous to (2), with $K^0$ replacing the nucleus as source, $\chi$ exchange replacing $\gamma$ exchange, and the nucleus replacing the electron as the spinning object. Mathematically, we must substitute $g^2_\chi$ for $\sigma$, $2K_\chi$ for $1 + 2K$, nucleon mass for electron mass, and $K$ velocity in the laboratory (or equivalently nuclear velocity in the $K$ rest frame) for electron velocity. We find

$$V_{\chi_{\text{exchange}}} = -(K_\chi g^2_\chi / 2m_\chi s t^2) \sigma \cdot r \times v.$$
together, the spin-orbit terms cancel unless the nuclear spins are polarized. Under standard operating conditions for the bubble chamber, for instance at $25^\circ\text{K}$ in a magnetic field of $10^4$ G, the excess of nuclear spins along the field is about one in $10^3$. As a result, only about $10^{-6}$ of the nuclei act together to produce a coherent $\chi$-exchange potential. From the directional properties of $\mathbf{a} \times \mathbf{r} \times \mathbf{v}$, one sees that the coherent potential is greatest for $K$'s passing through the chamber off-center, and vanishes when there are as many nuclei to the left as to the right.

Let us estimate the sum of $\mathbf{a} \times \mathbf{r} \times \mathbf{v}$ terms for a $K_2$ with $v \approx c$ approaching a bubble chamber a few centimeters off center. We suppose the bubble chamber has a (10 cm$^3$) volume, $\sim 10^{20}$ nuclei/cm$^3$, and in one of the nuclear spins aligned, and we give $K_2$ the small value 0.01 to avoid any conflict with the limit imposed by Ramsey's experiment (Sec. II-B). The result is

Total $V_{\chi\text{-exchange}} \sim (10^{17} \text{ nuclei})$

\[
\times (10^{-5} \text{ fraction of aligned spins})
\times (10^{-24} \text{ eV/nucleus}) \sim 10^{-2} \text{ eV}.
\]

(4)

Since this off-diagonal matrix element connecting $K_2$ to $K_1$ is much greater than the diagonal matrix element distinguishing $K_1$ from $K_2$ (approximately $\sim 4 \times 10^{-8}$ eV), the eigenvectors of $K_0$ are each $50\%$ of $K_2$ and $50\%$ of $K_1$, and the mass eigenvalues are split by $2V \sim 2 \times 10^{-24}$ eV. An initially pure $K_2$ beam would thus oscillate rapidly between $K_2$ and $K_1$, and many $K_1$ decays would occur as the beam approached and entered the chamber. This behavior is not observed, so again we must reject the possibility that $m_\chi$ is very small.

Now as $m_\chi$ is increased, several factors act to cut down $K_1$ regeneration: the $K_2$ beam doesn't exchange $\chi$'s with the bubble chamber until it gets closer than $r \approx 1/m_\chi$; only the nearby parts of the chamber, amounting to a volume $\sim m_\chi^{-3}$, contribute to the potential when $1/m_\chi \lesssim 10$ cm, and the left- and right-hand $\mathbf{a} \times \mathbf{r} \times \mathbf{v}$ contributions cancel for $K_2$ inside the chamber except in a peripheral region of thickness $1/m_\chi$. Taking these factors into account, we estimate that masses $m_\chi \lesssim (1 \text{ cm})^{-1}$ produce too much $K_1$ regeneration and can be ruled out.

III. ANOMALOUS $K_1$ REGENERATION

In a recent experiment, Leipuner et al.\textsuperscript{9} claim to observe anomalously large $K_1$ regeneration from a $K_2$ beam in a hydrogen bubble chamber. The anomalous $K_1$'s appear at very small angles with $\cos \theta \geq 0.999$, indicating that the momentum transfer $|t|$ to the hydrogen target is less than $(30 \text{ MeV})^2$. Other studies\textsuperscript{10} at $1/m_\chi \sim 1$ cm, only about $10^{-3}$ of the nuclei in the chamber are near enough to exchange $\chi$'s, so the total potential is reduced to $\sim 10^{-6}$ eV. This is already slightly smaller than the $K_1-K_2$ mass difference and thus only a fraction of the $K_2$'s exposed to such a potential will convert to $K_1$. Furthermore, the maximum potential will not be felt by $K_1$'s entering the central parts of the chamber, because of the left-right cancellation. Further increase of $m_\chi$ reduces the $K_1$ production very rapidly.

of $K_1$ regeneration in iron\textsuperscript{18} show that the anomaly cannot increase strongly when heavier nuclei are used as targets.

If we did not have the stringent limitations on $m_\chi$ recorded in Sec. II and IV, it would be possible to explain the results of Leipuner et al., in terms of $\chi$-exchange contributions to the reaction $K_2^+ (\text{nucleus}) \rightarrow K_1^+ (\text{nucleus})$. The $\chi$ would need to have $m_\chi < 30 \text{ MeV}$ to explain why the effect is concentrated at small angles. The failure of the effect to increase strongly in heavier nuclei would be explained naturally since $\chi$-exchange couples to nuclei only through the "strange" magnetic moment, which cancels in closed shells and does not increase proportionately to the number of nucleons.

The result of Leipuner et al. could then emerge in either of two ways. One possibility is $K_1$ regeneration in the collective field of many nuclei, as described in Sec. II-C, with $m_\chi$ just on the borderline between too much $K_1$ regeneration and too little [i.e., for the choice of $K_2$ in Sec. II-C, we would have $m_\chi \sim (1 \text{ cm})^{-1}$]. In this case, however, most of the $K_1$'s would be produced at the extreme sides of the chamber, where the $\mathbf{a} \times \mathbf{r} \times \mathbf{v}$ contributions from different nuclei do not cancel.

Another possibility would involve larger $m_\chi$ for which the collective exchange potential is negligible. The $K_1$-regeneration potential would then become sizable only in the immediate neighborhood of individual nuclei [remember, Eq. (3) grows like $r^{-2}$], and $K_1$ regeneration would depend on the occasional encounters of $K_2$ with such regions, occurring with equal probability over the entire bubble chamber. In a hydrogen chamber, this effect can be described in terms of the cross section for $K_2^+ p \rightarrow K_1^+ p$. The cross section is easily estimated from the Rosenbluth formula by changing $\epsilon$ to $\epsilon_\chi$ in the vertex, $t^{-1}$ to $(l-m_\chi^{-2})^{-1}$, and keeping only the anomalous magnetic moment $K_1 \delta g_\chi/2m_\chi$. One finds

$$d\sigma/d\Omega \approx g_\chi K_1^2 \delta g_\chi/M^2 (\epsilon\Omega + m_\chi^{-2})^3$$

(5)

where $q$ is the $K$ momentum and $M$ the proton mass.\textsuperscript{16}

At the small angles where most of the cross section is concentrated, $\epsilon \approx 1 - q^2/M^2$, so

$$d\sigma/d\Omega \approx g_\chi K_1^2 \delta g_\chi/M^2 (\epsilon\Omega + m_\chi^{-2})^3$$

(6)

The integrated cross section is approximately

$$\sigma_1 \equiv \int d\Omega \frac{2\pi g_\chi K_1^2}{d\Omega} \frac{\ln \left( \frac{2q}{m_\chi^2} \right)}{M^2}$$

\[
\sim \frac{K_1^2}{4} \ln \left( \frac{2q}{m_\chi^2} \right) 10^{-48} \text{ cm}^2.
\]

(7)


\textsuperscript{16} Incidentally, the forward differential cross section does not become infinite when $m_\chi = 0$ because the $K_1-K_2$ mass difference prevents momentum transfer $t$ from approaching zero. This effect has been left out of Eqs. (6) and (7).
Equating $\sigma_f$ with the anomalous effect of Leipuner et al., with $m_\gamma$ in the relevant range $2 \times 10^{-6}$ eV < $m_\gamma$ < $3 \times 10^6$ eV, we find we must take $K_\chi$ between $\frac{1}{4}$ and 2.

To summarize, it is possible to explain the experiment of Leipuner et al. by $\chi$ exchange. The explanation is untenable, however, because it contradicts other experimental evidence (Table I) against small $m_\chi$, and if $m_\chi$ < $10^8$ eV, it also contradicts the Ramsey limit $K_\chi$ < 0.03.

IV. EVIDENCE AGAINST $m_\chi$ < $2m_e$

A. Vacuum Polarization

In ordinary vacuum polarization, the largest contribution comes from the process (virtual $\gamma$) $\rightarrow$ $e^+$ + $e^-$ (virtual $\gamma$). This process can also be thought of as providing an electron-pair exchange correction to the one-photon exchange potential between two particles. For example, the Coulomb potential is corrected as follows:

$$ V = \frac{Q_1 Q_2}{r} + \frac{Q_1 Q_2}{2 m_e} \int \frac{dm' C(m') e^{-m' r}}{r}, \quad (8) $$

where the extra factor $\alpha$ in the correction term, and its short range $r$ ~ $1/2m_e$, are exhibited explicitly.

The $\chi$ is a spin-one particle and is negative under charge conjugation, just like the photon. Vacuum polarization can therefore receive a contribution from (virtual $\gamma$) $\rightarrow$ $\chi$ (for example, via $K^+K^-$, which couples to both $\gamma$ and $\chi$) $\rightarrow$ (virtual $\gamma$). This new process can be thought of as providing a $\chi$-exchange correction to the one-photon exchange potential (it also provides a $K^+K^-$ exchange potential, for example, but this has a very short range in an on-the-mass-shell theory and can safely be neglected). Thus to the Coulomb potential between two charged particles, one must add:

$$ \delta V = \frac{Q_1 Q_2 e^{-x}}{r} \frac{C_x}{x} \left( e^{-x r} / r \right). \quad (9) $$

We have estimated the residue of the $\chi$-exchange pole, for the purpose of obtaining $C_x$, by ordinary techniques of field theory. In our estimate, $\chi$ is joined to the charge line at each vertex by $\chi \rightarrow K^+K^- \rightarrow \gamma \rightarrow$ (charge line). After current conservation is imposed, the calculation is still logarithmically divergent so we introduce a cut-off mass $\Lambda$. The result for $m_\chi$ < $m_K$ is

$$ C_x \approx \frac{1}{36\pi^2} (\ln \Lambda / m_K)^2. \quad (10) $$

It is now clear that $\delta V$ is not likely to be more than a $10^{-4}$ correction to the Coulomb potential, so its effects will be noticeable only in relatively precise measurements. The Lamb shift provides a useful example. The potential $\delta V$ will shift an atomic level with wave function $\psi$ by approximately

$$ \Delta E = \int \psi^* V \psi d^3r. \quad (11) $$

For example, the 2P level in hydrogen undergoes the shift

$$ \Delta E(2P) = \int_0^\infty r^2 dr \left( -\frac{\varepsilon \alpha \alpha g^2 C_x e^{-m r}}{r} \right) \left( \frac{r e^{-r/a}}{(2a)^{3/2} \sqrt{3a}} \right)^2 $$

$$ = \frac{\varepsilon \alpha \alpha g^2 C_x}{4a (1 + a m_x)^4}, \quad (12) $$

where $a$ is the Bohr radius. Similarly, the 2S level in hydrogen undergoes the shift

$$ \Delta E(2S) = \left[-\frac{\varepsilon \alpha \alpha g^2 C_x}{4a (1 + a m_x)^4} \right] (1 + 2a^2 m_x^2). \quad (13) $$

The difference

$$ D = \Delta E(2P) - \Delta E(2S) = \varepsilon \alpha \alpha g^2 C_x a m_x^2 / (1 + a m_x)^4 \quad (14) $$

represents a correction to the Lamb shift, and must be bounded by 0.1 Mc/sec to maintain the agreement between experiment and standard quantum electrodynamics. If we take $\ln(\Lambda / m_K) = 2$ for purposes of estimating $C_x$, $D$ is large and the 0.1 Mc/sec agreement is very badly violated when $m_\chi$ lies between about 2 MeV and 8 eV, so we can rule out these values of $m_\chi$.

The behavior of $D$ can be given a simple qualitative explanation. We begin with the usual observation that $\delta V$ is most important at small $r$ ($r \leq 1/m_x$), where it influences the 2S more than the 2P level. As $m_\chi$ is decreased from high values, $\delta V$ overlaps an increasing area of the atom and thus $D$ grows until $\delta V$ covers the whole atom ($m_\chi \sim 1$). When $m_\chi$ is decreased beyond this point, most of the further growth of $\delta V$ occurs outside the atom, with the result that $\Delta E(2S)$ and $\Delta E(2P)$ approach constant values. The remaining changes of $\delta V$ inside the atom are towards a strictly 1/r behavior which shifts 2P as much as 2S, so $D$ slowly subsides back to zero.

As $m_\chi$ is reduced still farther to about $10^{-4}$ eV ($1/m_\chi \sim 1$ cm), the $\chi$-exchange potential conflicts with another very precise measurement: the Cavendish "ice-bucket" experiment which accurately verifies the 1/r form of the Coulomb potential. The result of the modern version of this experiment, performed by Plimpton and Lawton, is usually stated in terms of an $r^{-3/4}$ form for the Coulomb potential, and reads:

$$ |\delta| < 2 \times 10^{-9}. \quad (15) $$

For our purposes, we compare the experiment with the potential

$$ V = (Q_1 Q_2 / r) (1 + \alpha g^2 C_x e^{-m r}) \quad (16) $$

(ordinary vacuum polarization is entirely negligible at the distances, of order centimeters, involved in this experiment). The second term in parentheses is compatible with the Plimpton-Lawton experiment if it is

\[17\] If $\ln(\Lambda / m_K)$ exceeds two, more stringent limits on $m_\chi$ are obtained.

either very small at \( r \sim \text{centimeters} (m_x > 1/\text{cm}) \) or essentially constant over the apparatus \( (m_x \text{ very small}) \). Putting in numbers in a very crude fashion, one finds that the intermediate masses \( m_x \sim 10^{-13} \text{eV} (1/m_x \sim 1 \text{ cm}) \) to \( m_x \sim 10^{-10} \text{ eV} (1/m_x \sim 10^{-13} \text{ cm}) \) are incompatible with experiment.

B. \( \pi^0 \rightarrow \gamma + \chi \)

The decay \( \pi^0 \rightarrow 2\gamma \) is relatively slow for an electromagnetic process. Very likely, the explanation is that the decay has to proceed through rather massive intermediate states.\(^{20}\) In this case, intermediate states involving strange particle pairs may be competitive, and these pairs would couple to \( \chi \)'s more strongly than to \( \gamma \)'s, so we are led to consider the possibility of \( \pi^0 \) decay into \( \chi \)'s.

The decay \( \pi^0 \rightarrow 2\chi \) happens to be inhibited, since \( \pi^0 \) has isotropic spin one, \( \chi \) has isotropic spin zero, and the \( \chi \) coupling conserves isotropic spin. The electromagnetic interactions needed to violate isotopic spin bring a factor \( \alpha \) into each matrix element, more than offsetting the advantage gained by replacing double-photon emission \( \sim \gamma g_\chi \) in the matrix element) with double-\( \chi \) emission \( \sim g_\chi^2 \) in the matrix element), and the net result is the negligible branching ratio

\[
\text{Rate } (\pi^0 \rightarrow 2\chi) = \frac{\text{Rate } (\pi^0 \rightarrow 2\gamma)}{(g_\chi^2)}/100. \tag{17}
\]

In the decay \( \pi^0 \rightarrow \chi + \gamma \), however, the photon coupling supplies the necessary violation of isotopic spin, and if strange particle pairs are sufficiently well represented in intermediate states we may have a ratio as large as

\[
R = \frac{\text{Rate } (\pi^0 \rightarrow \chi + \gamma)}{\text{Rate } (\pi^0 \rightarrow 2\gamma)} \sim \frac{g_\chi^2}{\alpha} \sim 10. \tag{18}
\]

The admittedly uncertain theoretical estimate (18) can be compared with experiments which limit the branching ratio \( R \). The best limit\(^{19}\) is obtained from absolute rates of the various processes that occur when \( \pi^0 \)'s stop in hydrogen:

\[
\begin{align*}
\pi^0 + p &\rightarrow n + \pi^0, &\pi^0 + p &\rightarrow \gamma + \gamma, &\pi^0 + p &\rightarrow \gamma + e^+ + e^-, &\pi^0 + p &\rightarrow n + \gamma, \tag{a} \\
&\rightarrow n + \pi^0, &\pi^0 &\rightarrow \gamma + e^+ + e^-, &\rightarrow n + \gamma, &\rightarrow n + e^+ + e^-. &\tag{b} \\
&\rightarrow n + \pi^0, &\pi^0 &\rightarrow \gamma + e^+ + e^-, &\rightarrow n + \gamma, &\rightarrow n + e^+ + e^- &\tag{c} \\
&\rightarrow n + \pi^0, &\pi^0 &\rightarrow \gamma + e^+ + e^-, &\rightarrow n + \gamma, &\rightarrow n + e^+ + e^- &\tag{d}
\end{align*}
\]

Samios,\(^ {21}\) for example, reports measurements of the total number of \( \pi^0 \) stops \([ (a) + (b) + (c) + (d) + \text{negligible processes}] \), the total number of \( e^+ e^- \) pairs \([ (b) + (d)] \), and the ratio of \( (b) \) to \( (d) \) \([ (b) \text{ is distinguished from (d) by kinematical considerations}] \). Combining these measurements with our very accurate knowledge\(^ {22}\) of the ratios \( (a)/(b) \) and \( (c)/(d) \), one has a complete specification of \( (a) \) through \( (d) \) which should agree with the total number of \( \pi^0 \) stops—unless some fraction of \( \pi^0 \)'s decay into \( \gamma + \chi \). In this case, either \( m_x < 2m_x \), so that \( \chi \) cannot decay into an \( e^+ e^- \) pair and too few electron pairs would be seen, or \( m_x > 2m_x \), in which event \( \chi \rightarrow e^+ e^- \) via an intermediate photon is the main decay mode of \( \chi \) and too many \( e^+ e^- \) pairs would be seen (if \( m_x \) is a substantial fraction of \( m_x \), the \( e^+ e^- \) pairs would also have a different energy from the usual Dalitz pairs).

In Samios' study, the stopped \( \pi^0 \) are accurately accounted for by the individual processes \( (a), (b), \) \( (c), \) and \( (d) \), and one obtains the branching ratio \( R = 0.04 \pm 0.05 \) for decay into \( \gamma + \chi \). Actually, this 10% limit applies to the case \( m_x < 2m_x \); for \( m_x > 2m_x \), a better limit of order 0.1% is obtained since nearly all \( \chi \) would then decay into \( e^+ e^- \), whereas normal \( \pi^0 \rightarrow \gamma + \gamma \) decays convert into Dalitz pairs only about 1% of the time.

Comparing these limits on the \( \pi^0 \rightarrow \chi + \gamma \) decay with the theoretical estimate (18), we may conclude that \( m_x > m_x \) with some confidence even though the theoretical estimate was quite uncertain.

C. More on the Production of \( \chi \) and Detection of \( \chi \rightarrow e^+ e^- \)

In any process where electric charge is accelerated, photon emission occurs. Most of the photons have low energies, but with sizable accelerations one obtains energetic photons (more than 1 MeV, say) in \( \sim 1/\alpha \) of all events.

Analogously, in processes where strangeness is accelerated, \( \chi \) emission occurs, unless too little energy is available to produce the extra \( m_x e^2 \). In processes where strange particles are scattered or produced in pairs, the analogy to photon emission is very close and energetic \( \chi \)'s are typically emitted in \( g_\chi^2/\pi \sim 3\% \) of the events. In processes where strangeness is violated, such as \( K \) decay, the analogy is less close because of the absence of current conservation,\(^ {23}\) but the probability of energetic \( \chi \) emission is still expected to be of order 3% or greater.

Although \( \chi \)'s are so easily produced, we have thus far paid relatively little attention to how one might detect them directly. For \( \chi \) masses in the range \( 0 < m_x < 2m_x \), there were reasons for this neglect. The main decay mode is \( \chi \rightarrow 3\gamma \) (remember, \( \chi \) is negative under charge conjugation). This decay is slow and not conspicuous. If \( \chi \)'s are absorbed instead of decaying, the events might be mistaken for photon absorptions at first sight.

In the range \( 2m_x < m_x < 2m_x \), by contrast, \( \chi \rightarrow e^+ e^- \) is the main decay mode. It proceeds rapidly via an intermediate photon (the \( \chi \) coupling can be estimated.

\(^{19}\) We are greatly indebted to Dr. Hans Kohrak, who pointed out the significance of \( \pi^0 \) decays for our inquiry, and provided the information on the relevant experiments.

\(^{20}\) For a recent statement of this idea, see J. B. Bronzan and F. E. Low, Phys. Rev. Letters 12, 522 (1964).


\(^{23}\) Note added in proof. For a study of this case, see S. Weinberg, Phys. Rev. Letters 13, 495 (1964).
using an intermediate $K^+K^-$ state as in Sec. IV-A), and shows up as a conspicuous electron pair at or near the point of $\chi$ production. Thus, to take $K_1$ decays as an example, $3\%$ or more of all decays would involve $\chi's:\ K^+ \rightarrow \pi^+ + \pi^- + \chi, \ K^0 \rightarrow \pi^+ + \pi^- + \chi,$ and so forth—and the $\chi's$ would make easily detectable $\pi^+ \pi^-$ pairs at the point of $K_1$ decay. One can put experimental limits of a few tenths of a percent or better on the fraction of such decays.\footnote{A discussion of the experimental situation with Dr. J. van Putten was very helpful.} however, allowing us to rule out $m_\chi$ in the range $2m_\pi < m_\chi < 2m_\tau.$ At higher $\chi$ masses, this method is not applicable because the main decay mode becomes $\chi \rightarrow 2\pi$ which is not so conspicuous.

V. COMMENTS ON MASSIVE $\chi's$

There are many other effects of light or massless $\chi's$ which would be harder to observe than the cases we have mentioned, but nevertheless interesting. For example, the reader may have noticed that exchange of massless $\chi's$ would produce a Coulomb-like attraction between particles of opposite strangeness such as $K^-K^+$ or $\Sigma^-K^+$, leading to Bohr-type bound states. But we hope everyone is convinced by now that $m_\chi > 2m_\pi,$ and we proceed to a brief discussion of higher $\chi$ masses.

A massive $\chi$ would appear as an $I=0, J^P=1^-$ resonance which is produced and decays, mainly into strongly interacting particles, at a rate intermediate between the rates characteristic of strong and electromagnetic interactions. Since $\chi$ couples directly to the strangeness current, decays to strange particle pairs would be somewhat preferred if $m_\chi > 2m_\pi.$

Ne'eman\footnote{Note added in proof. For the particular decay $K^+ \rightarrow \pi^+ + e^+ + e^-$, the very low branching ratio
$$\Gamma(K^+ \rightarrow \pi^+ + e^+ + e^-) \lesssim 1.1 \times 10^{-4}$$
has now been established [U. Camerini, D. Cline, W. F. Fry, and W. M. Powell, Phys. Rev. Letters 13, 318 (1964)].} has suggested the $\phi$ as a possible candidate, in view of its preference for $K\bar{K}$ decays and rather small decay rate.\footnote{In Ref. 3, we dissociated the $\chi$-strange current interaction from the additional possibility, mentioned by Ne'eman, of a because its interactions are quite strong. It is also conceivable that $\chi$ is lighter than these well-known vector mesons, but has been missed in the past because of its relatively weak couplings.} The $\omega$ is not such a good candidate.

We close with a comment concerning the current that $\chi$ couples to. In general, isotopic spin and hypercharge conservation only allow $\chi$ to couple to a linear combination of the baryon ($B$) and hypercharge ($Y$) currents. The reason for taking the linear combination to be precisely the strangeness current $S = Y - B$ was to make a light or massless $\chi$ hard to observe—had it coupled to nonstrange currents as well, it would have been too easily produced. Our finding that $m_\chi$ must exceed $2m_\pi$ removes this motivation, allowing $\chi$ to couple to a more general linear combination of $Y$ and $B$ currents.\footnote{In particular, it is conceivable that $\chi$ could have a strong coupling to the $B$ (singlet) current, and substantially weaker coupling to the $Y$ (octet) current. The current-current product which appears in lowest order ($g_\chi^2$) mass splittings would then have a strong 1 component, a weaker 8 component, and a still weaker 27 component, which would help to explain why SU(3) symmetry violation is mainly octet. In view of its strong coupling to the $B$ current, $\chi$ could then be $\omega.$ This theory, however, goes beyond Ne'eman's original intentions inasmuch as it uses $\chi$ coupling to explain several things [strong singlet vector meson-baryon current coupling, octet dominance of SU(3) violation] which the bootstrap mechanism may be able to account for without introduction of a new interaction.}

ACKNOWLEDGMENTS

We would like to thank W. Chinowsky, H. Kobrak, Y. Ne'eman, J. van Putten, and W. Wagner for very helpful discussions.

$\chi$-muon current interaction. This additional possibility can indeed be ruled out if $\phi$ is $\chi.$ The decay width for $\phi \rightarrow \mu^+ + \mu^-$ would be
$$\Gamma = g_\chi^2 (m_\mu/3)[1 + O(m_\mu/m_\phi)] \sim 30 \text{ MeV}$$
if the muon had strangeness, but the experimental full width for $\phi$ is only 3 MeV \cite{N. Gelfand, D. Miller, M. Nussbaum, J. Ratna, J. Schultz et al., Phys. Rev. Letters 11, 438 (1963),} and the mode $\phi \rightarrow \mu^+ + \mu^-$ occurs in less than 1.3\% of $\phi$ decays. \cite{A. Barbaro-Galtieri and R. Tripp, University of California Report No. UCRG-11428, 1964 (unpublished).} This work will appear in the Proceedings of the International Conference on High Energy Nuclear Physics, Dubna, USSR, 1964 (to be published). We would like to thank Dr. Barbaro-Galtieri for informing us of her work in advance of publication.}