Higgs-meson production in muon-number-violating processes

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We calculate cross sections for coherent production of Higgs mesons in the Coulomb field of a heavy nucleus in processes which violate conservation of muon number. Depending on the specific model for muon-number violation and on the mass of the Higgs meson, such processes could be observable at Fermilab.

I. INTRODUCTION

Experiments on nuclear double-β decay provide overwhelming confirmation of the law of electronic-lepton-number conservation for energies below the threshold for muon production,[1] and a variety of experiments in which muons are energetically allowed strongly supports the existence of an additional and separate conservation law for muonic lepton number.[2] Nevertheless, violation of two separate conservation laws at the superweak level has been the subject of speculation for some time,[3] and such a violation appears more often than not in modern gauge theories.[4] Consequently, experimental efforts to observe violations of separate electronic- and muonic-lepton-number conservation laws have been undertaken with renewed vigor.[5]

We consider here the possibility of observing muon-number violation by coherent production of a Higgs boson in the Coulomb field of a heavy nucleus. As a specific theoretical framework for the calculation we start with the model recently proposed by Bjorken and Weinberg.[6] We first calculate the processes

$$\mu + A \to e + \phi^0 + A\,, \tag{1a}$$

$$\nu_\mu + A \to e + \phi^+ + A \tag{1b}$$

in the Coulomb field of a heavy nucleus, A. These processes occur through the couplings $\bar{e}\mu\phi$ and $\bar{e}\nu_\mu\phi$, which have the strength $m_\mu G_F^{1/2}$ in the model. Secondly, we consider a possible extension of the Bjorken-Weinberg model, which includes the charged heavy lepton U of Perl *et al.*,[7] in addition to ordinary leptons. The extended model contains a $\phi U \nu_\mu$ (ϕ = Higgs meson) vertex, which is considerably stronger than the analogous $\phi e \nu_\mu$ and $\phi \mu \nu_e$ vertices of the original model, because the coupling is proportional to m_U. (Because of the large coupling, care must be taken to arrange the extension so that the model does not violate the experimental upper limit for $\mu \to e\gamma$.) Finally, we calculate the cross section for the coherent process

$$\nu_\mu + A \to U^- + \phi^+ + A\,. \tag{2}$$

II. HIGGS PRODUCTION IN THE ORIGINAL MODEL

In the SU(2) × U(1) gauge model of Bjorken and Weinberg,[6] separate electronic- and muonic-lepton-number conservation is not violated by the gauge-field interactions. The violation occurs through the interaction of a number of Higgs doublets with leptonic matter according to

$$H_I = +g_1(\bar{\nu}_\mu, \overline{\mu^-})_L \begin{pmatrix}\phi_1^+\\ \phi_1^0\end{pmatrix}\mu_R^- + g_2(\bar{\nu}_e, \overline{e^-})_L\begin{pmatrix}\phi_2^+\\ \phi_2^0\end{pmatrix}\mu_R^- + g_3(\bar{\nu}_\mu, \overline{\mu^-})_L\begin{pmatrix}\phi_3^+\\ \phi_3^0\end{pmatrix}e_R^- + g_4(\bar{\nu}_e, \overline{e^-})_L\begin{pmatrix}\phi_4^+\\ \phi_4^0\end{pmatrix}e_R^- + \text{H.c.}\,, \tag{3}$$

where R and L stand for multiplicative factors of $\frac{1}{2}(1-\gamma_5)$ and $\frac{1}{2}(1+\gamma_5)$, respectively. The leptonic basis chosen for diagonalization of the mass matrix is such that

$$g_1\langle\phi_1^0\rangle = m_\mu\,,\quad g_2\langle\phi_2^0\rangle = g_3\langle\phi_3^0\rangle = 0\,,\quad g_4\langle\phi_4^0\rangle = m_e \tag{4}$$

and the very plausible assumptions are made that $\langle\phi_1^0\rangle$ and $\langle\phi_4^0\rangle$ are of order $G_F^{-1/2}$ and that at least one of g_2 and g_3 is comparable to g_1.

The model is further restricted by working in the $m_e = 0$ limit for which either $g_2 = g_4 = 0$ or $g_3 = g_4 = 0$. The calculations of Ref. 6 are done for the latter choice. While there is no substantial difference in their results if one makes the complementary choice $g_2 = g_4 = 0$, a nonzero value for g_3 provides a mechanism for both $\mu \to e\phi^0$ and $\nu_\mu \to e\phi^+$ in a nuclear Coulomb field. We shall

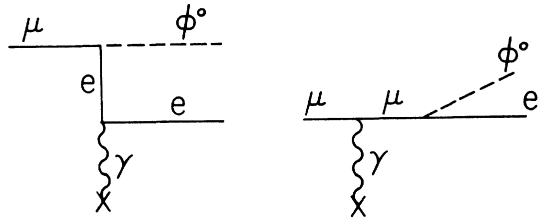

FIG. 1. Feynman diagrams for coherent production of a Higgs boson through $\mu \rightarrow e$ conversion in the Coulomb field of a nucleus.

consider both of these conversion processes.

The Feynman diagrams for processes (1a) and (1b) are shown in Figs. 1 and 2. The coherent production in a nuclear Coulomb field is mediated by one-photon exchange with the nucleus. The cross section for a reaction of this type can be adequately approximated by[8]

$$\frac{d^2\sigma}{dq^2dW^2}=\frac{Z^2\alpha}{\pi}\,|F(q^2)|^2\,\frac{q^2-\eta(W^2)}{q^4}\,\frac{\sigma_\gamma(W^2,q^2)}{W^2}\,, \qquad (5)$$

where

$$\eta(W^2)=q_{\min}{}^2=\frac{W^4}{4E^2}\,.$$

For reaction (1a) or (1b), $\sigma_\gamma(W^2,q^2)$ is the cross section for the related process $\mu+\gamma\rightarrow\phi^0+e$ or $\nu_\mu+\gamma\rightarrow\phi^++e$, respectively. For numerical purposes we use the exponential nuclear form factor

$$F(q^2)=\exp(-q^2R^2/6)\,,$$

with

$$R(\tfrac{3}{5})^{1/2}A^{1/3}/m_\pi\,. \qquad (6)$$

The virtual-photon mass is $-q^2$, W is the energy of $\phi+e$ in their c.m. frame, and E is the laboratory energy of the incident lepton. As $E\rightarrow\infty$, $\eta(W^2)\rightarrow 0$ so that there exists a kinematical region for which the form factor is unity, and the important q^2 region is limited by the nuclear charge radius R to lie between zero and $\sim 3/R^2$.

Except in the exchanged-lepton propagators, which become infinite when all masses vanish, we have let $\sigma_\gamma(W^2,q^2)\rightarrow\sigma_\gamma(W^2,0)$ and ignored lepton masses. In order to estimate the total cross sections for reactions (1a) and (1b), we have performed a numerical double integration of Eq. (5). The resulting cross sections on a $^{208}_{82}\mathrm{Pb}$ target, $\sigma(E)$, are shown for various possible $M=M_\phi$ in Figs. 3(a) and 3(b). Values of M significantly lower than those shown would lead to a rate for $\mu\rightarrow e\gamma$ in excess of the present experimental limit.[6]

These cross sections are rather small up to the highest Fermilab energies ($\leq 10^{-40}$ cm^2/nucleon), but leptonic decays of ϕ would provide a distinctive signal. If $g_2\neq 0$, one only has the muon-induced process, and its signatures are

$$\mu+A\rightarrow\begin{cases} e+e+\bar{\mu}+A\,, & (7a)\\ e+\mu+\bar{\mu}+A\,. & (7b)\end{cases}$$

The first of these requires ϕ_2^0 to be self-adjoint and the second requires a modest ϕ_1-ϕ_2 mixing. If $g_3\neq 0$, both μ- and ν_μ-induced processes are possible. The same signatures are still available for the muon process, and the neutrino process can be identified through

$$\nu_\mu+A\rightarrow e+\nu_\mu+\bar{\mu}+A\,. \qquad (8)$$

We close this section with a comment on the prospects for detection of the Higgs-meson track. The leptonic decay width for $\phi^+\rightarrow\bar{\mu}+\nu_\mu$ is given by

$$\Gamma(\phi\rightarrow\mu\nu_\mu)=G_F\,\frac{m_\mu{}^2}{M}\,\frac{(M^2-m_\mu{}^2)^2}{M^2+m_\mu^2} \qquad (9)$$

and for $M=4$ GeV corresponds to a partial lifetime of $c\tau=2\times10^{-6}$ cm, so that direct observation of the Higgs-meson track is impossible. The Higgs-meson coupling to quarks considered in Ref. 6 contributes a comparable partial width, since the muon mass in the coupling is replaced by the bare-quark mass.

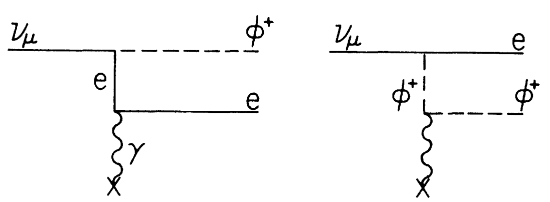

FIG. 2. Feynman diagrams for coherent production of a Higgs boson through $\nu_\mu\rightarrow e$ conversion in the Coulomb field of a nucleus.

III. EXTENSION TO HEAVY LEPTONS

In order to introduce the heavy lepton U on the same footing as e and μ, we replace the interaction of Eq. (3) by

$$H_I=\sum_{i,j=1}^{3} g_{ij}(\bar{\nu}_i\,,\bar{l}_i)_L\begin{pmatrix}\phi^+_{ij}\\ \phi^0_{ij}\end{pmatrix} l^-_{jR}+\mathrm{H.c.}\,, \qquad (10)$$

where $i,j=1,2,3$ correspond to U^-, μ^-, and e^-, respectively, and $l_1^-=U^-$ etc. Diagonalizing the mass matrix as in Ref. 6 implies

$$g_{ij}\langle\phi^0_{ij}\rangle=\begin{cases} m_i\,, & i=j\,,\\ 0\,, & i\neq j\,.\end{cases} \qquad (11)$$

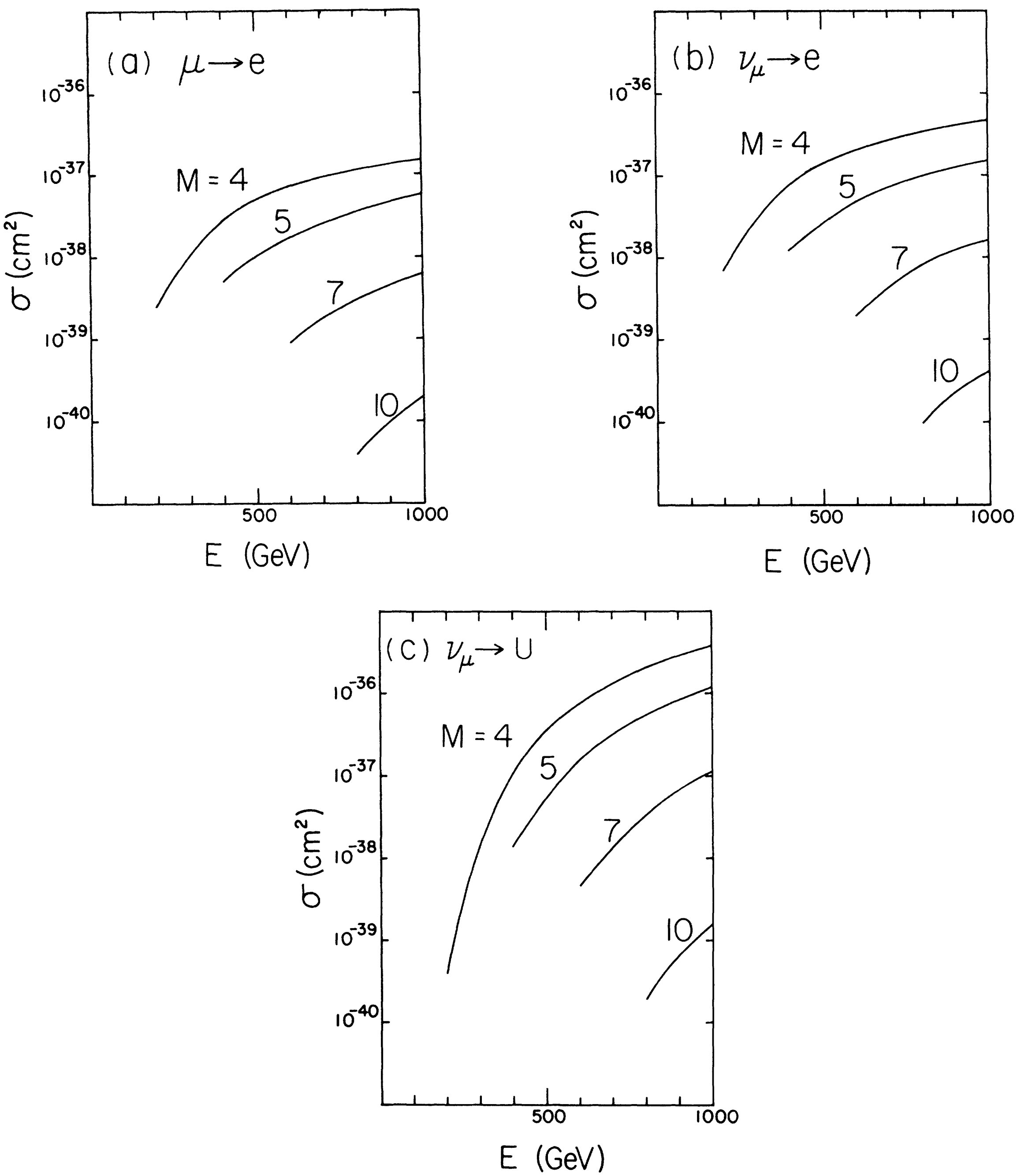

FIG. 3. Total coherent cross section for (a) $\mu + \mathrm{Pb} \to e + \phi^0 + \mathrm{Pb}$, (b) $\nu_\mu + \mathrm{Pb} \to e + \phi^+$, (c) $\nu_\mu + \mathrm{Pb} \to U + \phi^+ + \mathrm{Pb}$. Curves are shown for Higgs-meson masses of $M = 4, 5, 7, 10$ GeV.

Assumptions analogous to those of Bjorken and Weinberg clearly lead to an interaction of the form

$$G_F m_U(\nu_{\mu L} U_R^- \phi_{21}^+ + \mathrm{H.c.}) . \qquad (12)$$

It is the presence of the heavy-lepton mass factor here, instead of the much smaller muon mass in Eq. (4), which makes this muon-number-violating interaction so much stronger than any present in the original Bjorken-Weinberg model. If no restriction is made on the coupling g_{ij}, however, the process $\mu \to e\gamma$ can occur in this extended model via a one-loop diagram involving an intermediate U at the rate given by Eq. (3) of Ref. 6:

$$\frac{\Gamma_{\mu\to e\gamma}}{\Gamma_{\mu\to e\nu\bar\nu}}=\frac{12\alpha}{\pi}\;\frac{m_U{}^6}{m_\mu{}^2M_\phi{}^4}\,. \tag{13}$$

This exceeds the experimental upper limit[2] of 2.2×10^{-8} for $M_\phi<300$ GeV.

This difficulty with a straightforward extension of the model of Bjorken and Weinberg to include heavy leptons can be overcome in a natural way by arranging for *both* electron and muon to remain massless within the model, which now generates only the mass of the heavy lepton.[9] This can be accomplished by demanding that both e_R and μ_R decouple; i.e., in Eq. (8) $g_{13}=g_{33}=g_{23}=g_{22}=g_{32}=g_{12}=0$. In this case the process $\mu\to e\gamma$ no longer occurs in the model. Higgs-meson production in a muon-number-violating process [Eq. (2)] still takes place, however (Fig. 2 with $e\to U$).

We have calculated the cross section for reaction (2) on a lead nucleus using the same method as for reaction (1). The results are shown in Fig. 3(c). As a consequence of the energy dependence of $q_{\min}$ and the form factor cutoff at large q^2, $\sigma(E)$ becomes asymptotic for $E\gg(M+m_U)^2R/\sqrt{12}$. For $M=4$ GeV, and $m_U=2$ GeV this requires $E\gg360$ GeV. Thus, threshold behavior is important at present Fermilab energies. In the asymptotic region, apart from logarithmic dependences, σ scales according to $\sigma(E)\sim(1/M^2)F(E/M)$ for $m_U/M\ll1$. It is clear from the figure that the cross section can become appreciable if M_ϕ is near 4 GeV but that if $M>5$ GeV, reaction (2) will be hard to detect (σ/nucleon $\lesssim10^{-40}$ cm^2) for $E_\nu\lesssim300$ GeV.

The reaction $\phi^+\to U^+\nu_\mu$ is expected to be an important decay mode. From the analog of Eq. (9), its rate for $m_U=2$ GeV and $M=4$ GeV corresponds to a partial lifetime of $c\tau=7\times10^{-9}$ cm, so that direct observation of the ϕ track is again impossible. A possible signature for detecting reaction (2) results when ϕ and U decay leptonically:

$$\begin{array}{l}\nu_\mu+A\to A+\phi^++U^-\\ \qquad\qquad\qquad\qquad\searrow\; l^-\bar\nu_l\nu_U\\ \qquad\qquad\searrow\\ \qquad\qquad\quad U^++\nu_\mu\\ \qquad\qquad\qquad\searrow\\ \qquad\qquad\qquad\; l^+\nu_l\bar\nu_U\,.\end{array} \tag{14}$$

This would lead to opposite-sign dilepton events ($\mu\mu$, μe, ee) with no hadronic energy. Existing data on dilepton events[10] could in principle set upper limits on reaction (14) and hence on (2). However, the cross section for (14) is proportional to B^2 (B = branching ratio for $U\to l\nu_e\nu_U$) and may therefore be suppressed considerably compared to the full cross section shown in Fig. 3(c).

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[4]For example, see F. Wilczek and A. Zee, Nucl. Phys. B106, 460 (1976); T. P. Cheng and F. Li, Phys. Rev. Lett. 38, 391 (1977).

[5]The $\mu\to e\gamma$ experiment underway at the Swiss Institute for Nuclear Research is well known by now, and we understand a new search for coherent neutrinoless muon capture will begin soon at the Los Alamos Meson Physics Facility.

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[8]A. Halprin, C. M. Andersen, and H. Primakoff, Phys. Rev. 152, 1295 (1966).

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