Revista Mexicana de Física **42**, *No. 4 (1996) 517–542*

Applications of effective Lagrangians to weak radiative decays

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Recibido el 11 de diciembre de 1995; aceptado el 22 de marzo de 1996

ABSTRACT. Using the effective Lagrangian approach we study the weak radiative processes $H^0 \to \gamma\gamma$, $\mu \to e\gamma$, $\tau \to \mu\gamma$, $e\gamma$ and $b \to s\gamma$. We assume that the full theory is perturbatively renormalizable and that nonrenormalizable interactions respects the symmetry of the standard model. We consider only those contributions generated at lowest level in the perturbative full theory. We find that the effects due to new physics may enhance the $H \to \gamma\gamma$ SM decay width and its production in $\gamma\gamma$ colliders up to one order of magnitude. We also get bounds on the fermion-number-nonconserving decays $\nu_i \to \nu_j\gamma$ and $t \to c\gamma$ as well on the anomalous electromagnetic couplings of the W boson and the t, c, u quarks.

RESUMEN. Se estudian los procesos radiativos débiles $H^0 \to \gamma\gamma$, $\mu \to e\gamma$, $\tau \to \mu\gamma$, $e\gamma$ y $b \to s\gamma$ mediante el esquema de lagrangianos efectivos. Se asume que la teoría completa es perturbativamente renormalizable y de que las interacciones no renormalizables respetan las simetrías del modelo estándar. Sólo se consideran contribuciones generadas al orden perturbativo más bajo en la teoría completa. Se encuentra que los efectos debidos a nueva física pueden incrementar la anchura del decaimiento $H^0 \to \gamma\gamma$ y su producción en colisionadores $\gamma\gamma$ hasta en un orden de magnitud con respecto a las predicciones del modelo estándar. También se obtienen cotas en los decaimientos

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que violan número fermiónico $\nu_i \to \nu_j\gamma$ y $t \to c\gamma$, así como en los acoplamientos electromagnéticos anómalos del bosón W y de los quarks t, c, u.

PACS: 12.15c; 14.60G; 14.65Ha

1. LINEAR EFFECTIVE LAGRANGIANS

The effective Lagrangian approach has been advocated as a means to parametrize physics beyond a given theory, especially within the context of the minimal standard model (SM) [1] and in the case that the low-energy scalar sector contains two doublets [2]. This approach parametrizes the low-energy limit of any Green's function in terms of the unknown coefficients of an, in principle, infinite set of effective operators involving only the low-energy fields. These operators respects the SM symmetries and are suppressed by inverse powers of the high-energy new physics scale Λ. The requirement that the low-energy particle content should be the same as in the SM is an assumption. While the requirement of model independence with respect to the full theory increases the number of unknown parameters whose order of magnitude can at best be estimated, only a finite number of them need to be considered in any given calculation; the number of operators which has to be considered is determined by the required degree of accuracy.

The results of the effective Lagrangian approach are very useful in determining, in a model-independent manner, the sensitivity of a given experiment to new physics, and can be used to isolate those observables most sensitive to possible new interactions. We assume that the underlying physics is described by a high-energy Lagrangian, out of which all heavy excitations with masses above ~ 400 GeV are integrated out; what remains will be the SM plus an infinite series of effective operators. There are two possible types of high-energy physics to consider: decoupling or non-decoupling from low-energy physics. In the weakly decoupled case, the masses of heavy degrees of freedom are large because a dimension full parameter respecting the symmetries of the theory is large. The decoupling theorem [3] then implies that all the effects of heavy excitations at low energies can be parametrized in a model-independent manner by a series of local operators which satisfy the same symmetries as the low-energy theory, and whose coefficients are suppressed by the appropriate power of Λ. In the nondecoupling case, the masses of the heavy degrees of freedom in the theory are large because some dimensionless coupling constant is large. In this scenario, the contributions due to physics above Λ need not be suppressed by inverse powers of Λ. The corrections to the low-energy processes are given by a chiral (nonlinear) expansion in powers of p/Λ, where p is a typical momentum for the given process [4]. In the present work we consider only the weakly coupled scenario since for strongly coupled underlying physics it is difficult to maintain the scalars naturally light as compared to Λ without considerably modifying the low-energy spectrum [5].

In the present paper we will review the possible effects of the physics underlying the SM in the weak radiative processes $H \to \gamma\gamma$, $b \to s\gamma$, $\mu \to e\gamma$, $\tau \to \mu\gamma$, $e\gamma$ and in the Higgs boson production in $\gamma\gamma$ colliders using the linear effective Lagrangian approach. We find that the effects due to new physics may enhance the SM $H \to \gamma\gamma$ decay width and the SM predictions for the Higgs boson production in the $\gamma\gamma$ fusion processes. We reexamine the possible deviations from the SM on the decays $b \to s\gamma$ and $\mu \to e\gamma$ within the context of the effective Lagrangian approach. Using the experimental constraints on these decay modes, we get bounds for the anomalous electromagnetic couplings of the W boson and t, c, u quarks,. as well as on the transition magnetic moments of the known neutrinos.

2. Two photon processes

In the minimal SM the CP-even $H\gamma\gamma$ interaction is generated through one-loop effects of charged fermions and W gauge bosons. The contribution arising from the W boson and the top quark are dominant, although the latter is marginally important and tends to cancel partially the first one [6]. It has been pointed out [7] that an intermediate-mass Higgs boson ($m_Z < m_H < 2m_Z$) could be detected through the rare decay mode $H \to \gamma\gamma$ at the LHC. It has been known for long that the $\gamma\gamma$ decay mode is rather sensitive to new physics effects [8]. Although the SM branching ratio $\mathrm{Br}(H \to \gamma\gamma) \sim 10^{-3}$ is small [6, 8], the importance of this decay mode lies in its sensitivity to gauge couplings and its clean signature.

The linear effective Lagrangian to be considered is thus given by

$$L_{\text{eff}} = L_0 + \sum_{n=5} \left[\sum_{i=1} \frac{\alpha_i}{\Lambda^{n-4}} (O_i^{(n)} + \text{H.C.}) \right], \tag{1}$$

where L_0 is the renormalizable SM Lagrangian and $O_i^{(n)}$ are nonrenormalizable operators of dimension $n \geq 5$ which respect the SM symmetries. Because of the absence of a $H\gamma\gamma$ tree-level interaction, the loop contributions must be finite. Any divergence can be associated with a local operator respecting the symmetries of the model; since the effective Lagrangian already contains all such operators, all divergences can be absorbed in a renormalization of the effective Lagrangian coefficients [1]. Thus, the cancellation of divergences, together with the constraint of electromagnetic gauge invariance, provide important checks on the algebra.

The unknown coefficients α_i in (1) depend on the underlying physics and are expected to be ≤ 1. If the underlying physics is weakly coupled, then the order of magnitude of the coefficients α_i is determined by the type of graph in the underlying theory which generates them. If an operator is generated at three level by the heavy dynamics, then α_i will be just a product of coupling constants of the underlying theory, which by assumption are ≤ 1 and thus also $\alpha_i \leq 1$. Loop generated operators will acquire an extra suppression factor $(1/16\pi^2)$. One should keep in mind that the results obtained with the effective Lagrangian (1) on low-energy processes represent thus an upper bound and could be suppressed significantly [9].

We will consider all the effects of new physics on the $H\gamma\gamma$ interaction which may be competitive with the SM (dimension-four) theory predictions [6]. Since this interaction occurs at the one-loop level in the SM, we consider only new effects up to the one-loop level in the full theory. Then, we have to consider three-level-generated dimension six fermionic operators which induce effective $H\bar{f}f$ interactions and contribute to $H\gamma\gamma$ through fermion loops [10]. There are not tree-level-generated dimension-six bosonic operators which could generate the $H\gamma\gamma$ interaction and we consider thus one-loop-generated dimension six ones. We will include also all the three-level-generated dimension-eight bosonic operators which lead to the $H\gamma\gamma$ interaction [11].

In Fig. 1 we display the Feynman diagrams which contribute to the $H^0 \to \gamma\gamma$ decay up to one-loop level in the full theory. There are six (CP-even) tree-level-generated dimension-six fermionic operators which contribute to this decay through the one-loop graphs shown in Figs. 1(e) and 1(f):

$$\begin{aligned} O_{\phi U} &= i(\phi^\dagger D_\mu \phi)(\bar{U}_{\rm R}\gamma^\mu U_{\rm R}), & O_{\phi D} &= i(\phi^\dagger D_\mu \phi)(\bar{D}_{\rm R}\gamma^\mu D_{\rm R}), \\ O_{\phi F}^{(1)} &= i(\phi^\dagger D_\mu \phi)(\bar{F}\gamma^\mu F), & O_{\phi F}^{(3)} &= i(\phi^\dagger D_\mu \tau^i \phi)(\bar{F}\gamma^\mu \tau^i F), \\ O_{U\phi} &= (\phi^\dagger \phi)(\bar{F} U_{\rm R}\tilde{\phi}), & O_{D\phi} &= (\phi^\dagger \phi)(\bar{F} D_{\rm R}\phi), \end{aligned} \tag{2}$$

where $U_{\rm R}$ and $D_{\rm R}$ denote right-handed fermions of type up and down, respectively, which are SU(2) singlets, and $\bar{F} = (\bar{U}, \bar{D})_{\rm L}$ denotes the corresponding left-handed SU(2) doublet; ϕ denote the scalar SU(2) doublet, τ^i are the Pauli matrices, $\tilde{\phi} = i\tau^2\phi^*$, and D_μ is the respective covariant derivative. There are three (CP-even) one-loop-generated dimension-six bosonic operators which induce tree-level effective $H^0\gamma\gamma$ and $H^0\gamma Z$ interactions:

$$\begin{gathered} O_{\phi W} = (\phi^\dagger \phi) W^i_{\mu\nu} W^{i\mu\nu}, \qquad O_{\phi B} = (\phi^\dagger \phi) B_{\mu\nu} B^{\mu\nu}, \\ O_{WB} = (\phi^\dagger \tau^i \phi) W^i_{\mu\nu} B^{\mu\nu}, \end{gathered} \tag{3}$$

where $W_{\mu\nu}$ and $B_{\mu\nu}$ are the SU(2) and U(1) field strength tensors, respectively.

The Higgs decay channel into $\gamma\gamma$ has been previously studied [12] using the CP-even operators. In order to get a prediction on the respective decay widths, these authors used the bounds on the coefficients $\varepsilon_i = \alpha_i(v/\Lambda)^2$ obtained from low-energy processes, which in general are of the order 10^{-2}. However, we consider that such bounds are still rather weak since in our scenario their contributions are expected to be suppressed because the respective coefficients α_i arise from dimension-six operators induced at the one-loop level of the complete theory [9]. These studies did not take into account the fact that effective interactions can be induced at different levels in the context of a simple perturbative scheme of the full theory.

Finally, there are five tree-level-generated dimension-eight bosonic operators which induce the $H^0\gamma\gamma$ interactions:

$$\begin{aligned} O_{8,1} &= (\phi^\dagger \tau^i \phi)(\phi^\dagger \tau^j \phi) W^i_{\mu\nu} W^{j\mu\nu}, & O_{8,2} &= i(D^\mu \phi^\dagger D^\nu \phi)(\phi^\dagger \tau^i \phi) W^i_{\mu\nu}, \\ O_{8,3} &= (\phi^\dagger \phi)^2 B_{\mu\nu} B^{\mu\nu}, & O_{8,4} &= i(\phi^\dagger \phi)(D^\mu \phi^\dagger D^\nu \phi) B_{\mu\nu}, \\ & & O_{8,5} &= (\phi^\dagger \phi)^2 W^i_{\mu\nu} W^{i\mu\nu} \end{aligned} \tag{4}$$

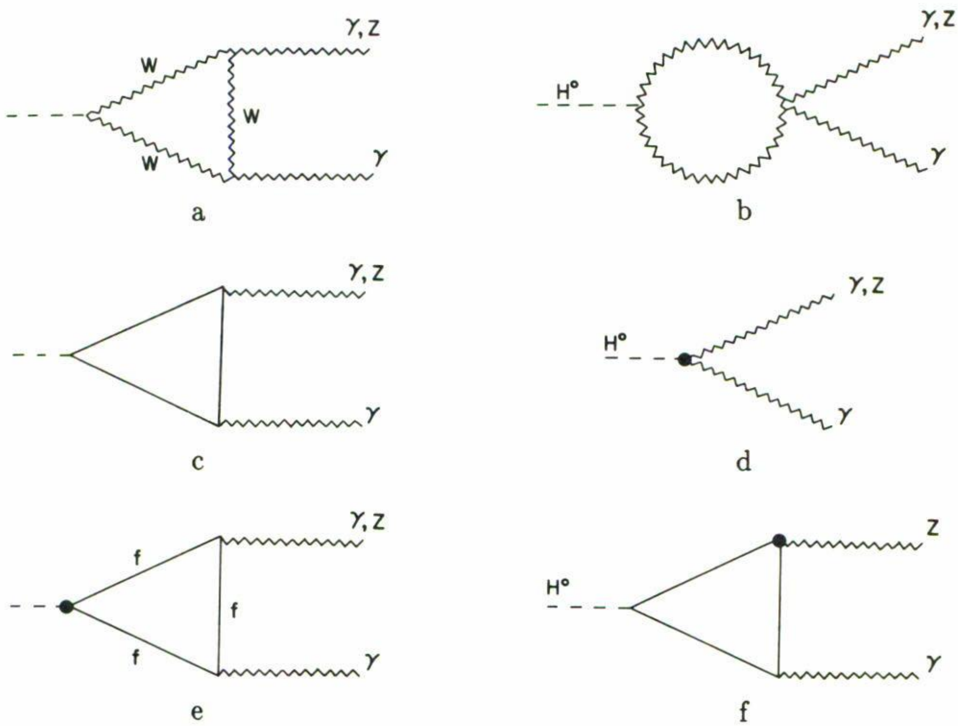

FIGURE 1. Feynman diagrams contributing to the $H^0 \to \gamma\gamma$, γZ decays. The heavy dot denotes an effective vertex. There are no three-level generated $\bar{f}f\gamma$ effective vertices.

Calculation of the Feynman diagrams shown in Fig. 1 leads to the following effective width for $H^0 \to \gamma\gamma$:

$$\Gamma_{\text{eff}}(H^0 \to \gamma\gamma) = a\left|\sum_f Q_f^2 N_C \left[F_0^f + \left(\frac{v}{\Lambda}\right)^2 F_a^f\right] + F_0^b + \left(\frac{v}{\Lambda}\right)^2 \left[F_{a6}^b + \left(\frac{v}{\Lambda}\right)^2 F_{a6}^b\right]\right|^2 \quad (5)$$

where

$$a = \frac{\alpha^2 G_{\text{F}} m_H^3}{64\sqrt{2}\pi^3}, \quad (6)$$

and the parametric functions F are given in Refs. [8, 11, 14].

In Ref. [11] we presented the ratio $\Gamma_{\text{eff}}/\Gamma_0$ as a function of m_H for $m_t = 174$ GeV, $\Lambda = 1$ TeV, and we took all the coefficients α_i equal to 1. In the fermionic part, only the quark top contributions were taken into account. We included the contributions coming from dimension-six fermionic operators, dimension-six bosonic operators, dimension-eight bosonic operators, and the overall contribution. It was found that this rare decay mode is insensitive to both fermionic and bosonic dimension-six contributions. In contrast, the dimension-eight tree-level-generated contributions enhance up to one order of magnitude the SM width $H^0 \to \gamma\gamma$ [11]. We have found that different scenarios for the values of the α_i coefficients do not modify appreciably this result. In Fig. 2 we present the same ratio

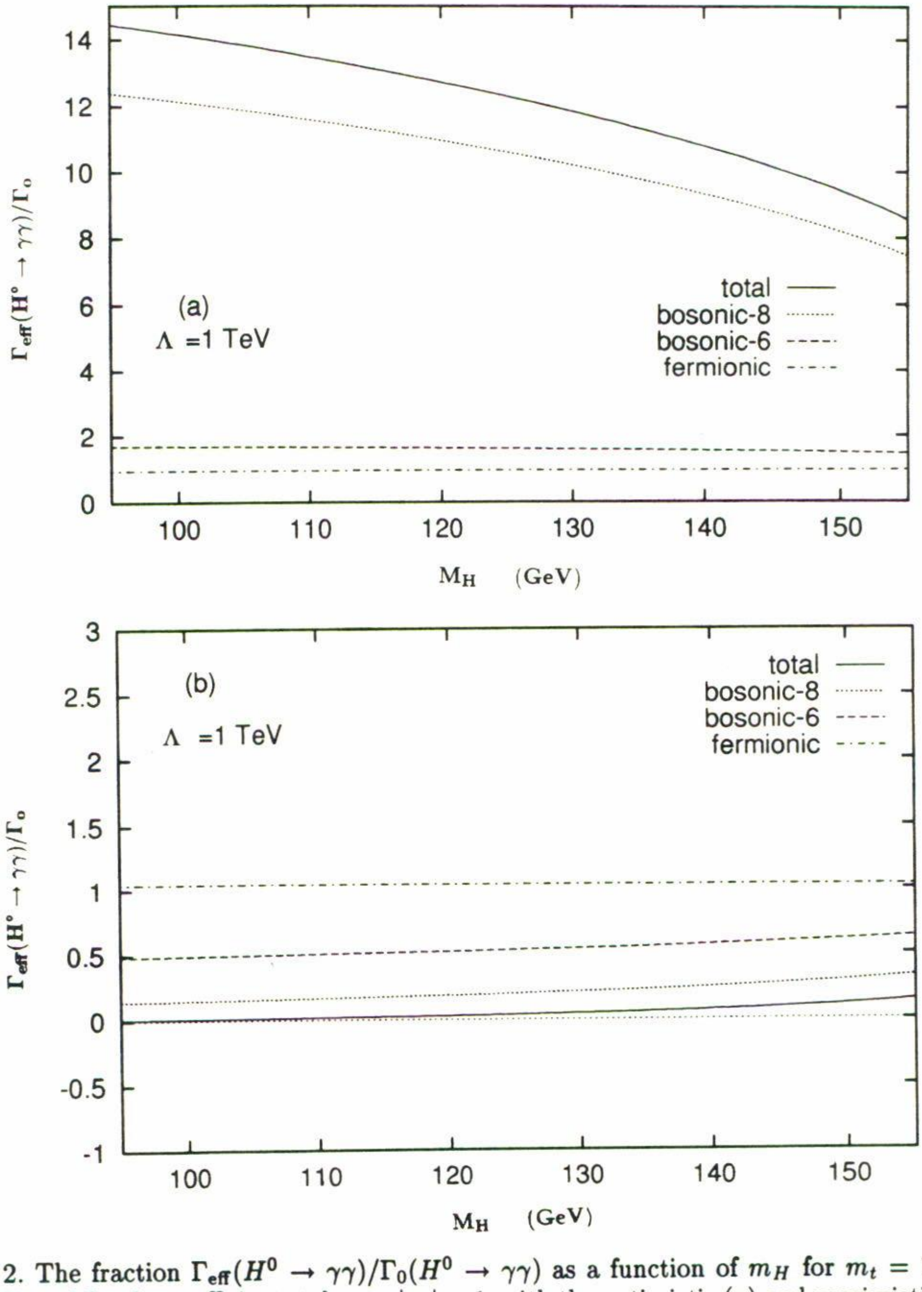

FIGURE 2. The fraction $\Gamma_{eff}(H^0 \to \gamma\gamma)/\Gamma_0(H^0 \to \gamma\gamma)$ as a function of m_H for $m_t = 174$ GeV, $\Lambda = 1$ TeV, and for the coefficients taken as $|\alpha_i| = 1$, with the optimistic (a) and pessimistic (b) scenarios. Fermionic (dot-dashed), dimension-six bosonic (dashed), dimension-eight bosonic (dotted), and total (solid line) contributions are displayed.

Γ_{eff}/Γ_0 but now we have displayed the optimistic and pessimistic scenarios taken with the constriction $|\alpha_i|$ equal to 1. Similar results were obtained for the decays $H \to \gamma Z$ and $Z \to \gamma H$ [11].

3. HIGGS BOSON PRODUCTION IN $\gamma\gamma$ COLLIDERS

The experimental progress in particle physics depends decisively on the set in operation of the CERN Large Hadron Collider (LHC) and the next generation of high energy e^+e^- linear colliders (NLC). An interesting option that is deserving a lot of attention nowadays is the possibility of transforming a linear e^+e^- collider in a $\gamma\gamma$ collider. By using the old idea of Compton laser backscattering it is possible to reach a center of mass energy of $\sqrt{s}$ = 200–500 GeV [13]. The $\gamma\gamma$ colliders will open a new window of research in experimental particle physics where the nature and source of the electroweak symmetry breaking may be tested. In particular, the $\gamma\gamma$ colliders offer a extraordinary opportunity to study the dynamics of the elusive Higgs boson. If the standard model (SM) Higgs is detected through its dominant decay modes in e^+e^- or $p\bar{p}$ collisions, then a $\gamma\gamma$ collider will allow a direct measurement of its partial decay width into two photons.

The Higgs production through $\gamma\gamma$ fusion has been recently studied in the minimal SM [2] and in the SM with two doublets [2]. In this section we present possible effects of new physics on the Higgs boson production in $\gamma\gamma$ fusion, within the context of the effective Lagrangian approach [1, 9, 11], in both the intermediate and heavy Higgs mass region, $m_W < m_{H^0} < 2m_W$ and $2m_W < m_{H^0} < 400$ GeV, respectively. We consider effective contributions in all the vertices involved taking as fundamental criterion the order in which the nonrenormalizable interactions may be generated by the heavy physics. We also present a qualitative discussion on the effects of new physics on the main sources of background expected in each mass region [14].

Since for $m_{H^0} \leq 400$ GeV the total width of the Higgs boson is small in comparison with the energy of the photons beam, the number of $H^0 \to X$ expected events may be written as [13]

$$N = \left[\frac{dL_{\gamma\gamma}}{dW_{\gamma\gamma}}\right]_{m_H} \frac{4\pi^2}{m_H^2}\Gamma(H^0 \to \gamma\gamma)\mathrm{BR}(H^0 \to X)(1+\lambda\tilde{\lambda}), \tag{7}$$

where $\mathrm{BR}(H^0 \to X)$ is the branching ratio of the Higgs boson into the final state X ($X = \bar{b}b$, WW, ZZ, $\bar{t}t$), λ and $\tilde{\lambda}$ are the helicity states of the scattered photons, $(dL_{\gamma\gamma}/dW_{\gamma\gamma})$ is the differential $\gamma\gamma$ luminosity as a function of the two-photon invariant mass. As it was already mentioned, we consider effects of new physics on the $H^0\gamma\gamma$, $H^0\bar{b}b$, $H^0\bar{t}t$, H^0WW and H^0ZZ interactions by taking into account the order in which they are generated by the underlying physics. Since the $H^0\gamma\gamma$ interaction is generated at one-loop level by the dimension-four theory, a complete calculation requires to consider all the non-renormalizable operators which contribute at this level in the full theory. This analysis includes operators of dimension six generated at tree-level which induce the $H^0\bar{f}f$ vertices and which contribute to $H^0\gamma\gamma$ through loops of fermions, and operators of dimension six, generated at one-loop level, which induces the $H^0\gamma\gamma$ interaction at tree-level. All these contributions to the decay $H^0 \to \gamma\gamma$, including the SM contribution, are suppressed by the loop factor $(16\pi^2)^{-1}$. Consequently, if the new physics scale Λ is not very far from the Fermi scale $v = \sqrt{2}\langle\phi\rangle_0$, the contributions of dimension eight operators, which may be generated at tree-level by the underlying physics, are dominant with respect to the contributions of the operators of dimension six [11]. Therefore, we will consider only the nonrenormalizable contributions induced by the operators of dimension eight which may be generated at tree-level. On the other hand, since the $H^0\bar{b}b$, $H^0\bar{t}t$, H^0WW and $H^\circ ZZ$ are tree-level interactions in the SM, the most important contributions of new physics to these vertices are given by those operators of dimension six which may be generated at tree-level by the underlying physics. Accordingly, the relevant effective Lagrangian for the Higgs boson production in $\gamma\gamma$ fusion may be written as

$$\begin{aligned}\mathcal{L}_{\text{eff}} = \mathcal{L}_0 &+ \frac{1}{\Lambda^2}\left[\alpha_{b\phi}O_{b\phi} + \alpha_{t\phi}O_{t\phi} + \alpha_\phi^{(1)}O_\phi^{(1)} + \alpha_\phi^{(3)}O_\phi^{(3)}\right] \\ &+ \frac{1}{\Lambda^4}\left[\alpha_{8,1}O_{8,1} + \alpha_{8,3}O_{8,3} + \alpha_{8,5}O_{8,5}\right],\end{aligned} \tag{8}$$

where $\mathcal{L}_0$ is the SM Lagrangian and α_i are unknown parameters; the operators of dimension eight and dimension six (2) and (4) induce the $H^0\gamma\gamma$ and $H^0q\bar{q}$ vertices at tree-level.

The operators of dimension six $O_\phi^{(1)}$ and $O_\phi^{(3)}$, which induce the H^0WW and H^0ZZ interactions at tree-level, and may be generated also at tree-level by the underlying physics, are given by

$$\begin{aligned}O_\phi^{(1)} &= (\phi^\dagger\phi)[(D_\mu\phi)^\dagger D^\mu\phi], \\ O_\phi^{(3)} &= (\phi^\dagger D_\mu\phi)[(D^\mu\phi)^\dagger\phi],\end{aligned} \tag{9}$$

where D_μ is the respective covariant derivative. Using these operators we obtain the following expressions for the Higgs partial decay widths:

$$\Gamma_{\text{eff}}(H^0 \to \gamma\gamma) = \frac{\alpha^2 G_F m_H^3}{64\sqrt{2}\,\pi^3}\left|\sum_f Q_f^2 N_c F_0^f + F_0^b + \left(\frac{v}{\Lambda}\right)^4 F_{a8}^b\right|^2, \tag{10}$$

$$\Gamma_{\text{eff}}(H^0 \to W^+W^-) = \left[1 + \frac{1}{4}\left(\frac{v}{\Lambda}\right)^2\left(2\alpha_\phi^{(1)} - \alpha_\phi^{(3)}\right)\right]^2 \Gamma_0(H^0 \to W^+W^-) \tag{11}$$

$$\Gamma_{\text{eff}}(H^0 \to ZZ) = \left[1 + \frac{1}{2}\left(\frac{v}{\Lambda}\right)^2\left(\alpha_\phi^{(1)} + \alpha_\phi^{(3)}\right)\right]^2 \Gamma_0(H^0 \to ZZ), \tag{12}$$

$$\Gamma_{\text{eff}}(H^0 \to \bar{q}q) = \left[1 - \frac{\sqrt{2}m_Z}{Cm_q}\left(\frac{v}{\Lambda}\right)^2\alpha_{q\phi}\right]^2 \Gamma_0(H^0 \to \bar{q}q), \tag{13}$$

where q stands for the b and t quarks, Q_f is the electric charge of the fermions, N_c is 3 for quarks and 1 for leptons, $C = m_Z\sqrt{\sqrt{2}\,G_{\text{F}}}$; and $\Gamma_0(H^0 \to W^+W^-)$, $\Gamma_0(H^0 \to ZZ)$, and $\Gamma_0(H^0 \to \bar{q}q)$ are the SM widths [11, 14]. The parametric functions F_0^i and F_8^i, associated to fermions and W boson loops, respectively, given in Refs. [11, 14], comprises the contribution of the operators (2), (4).

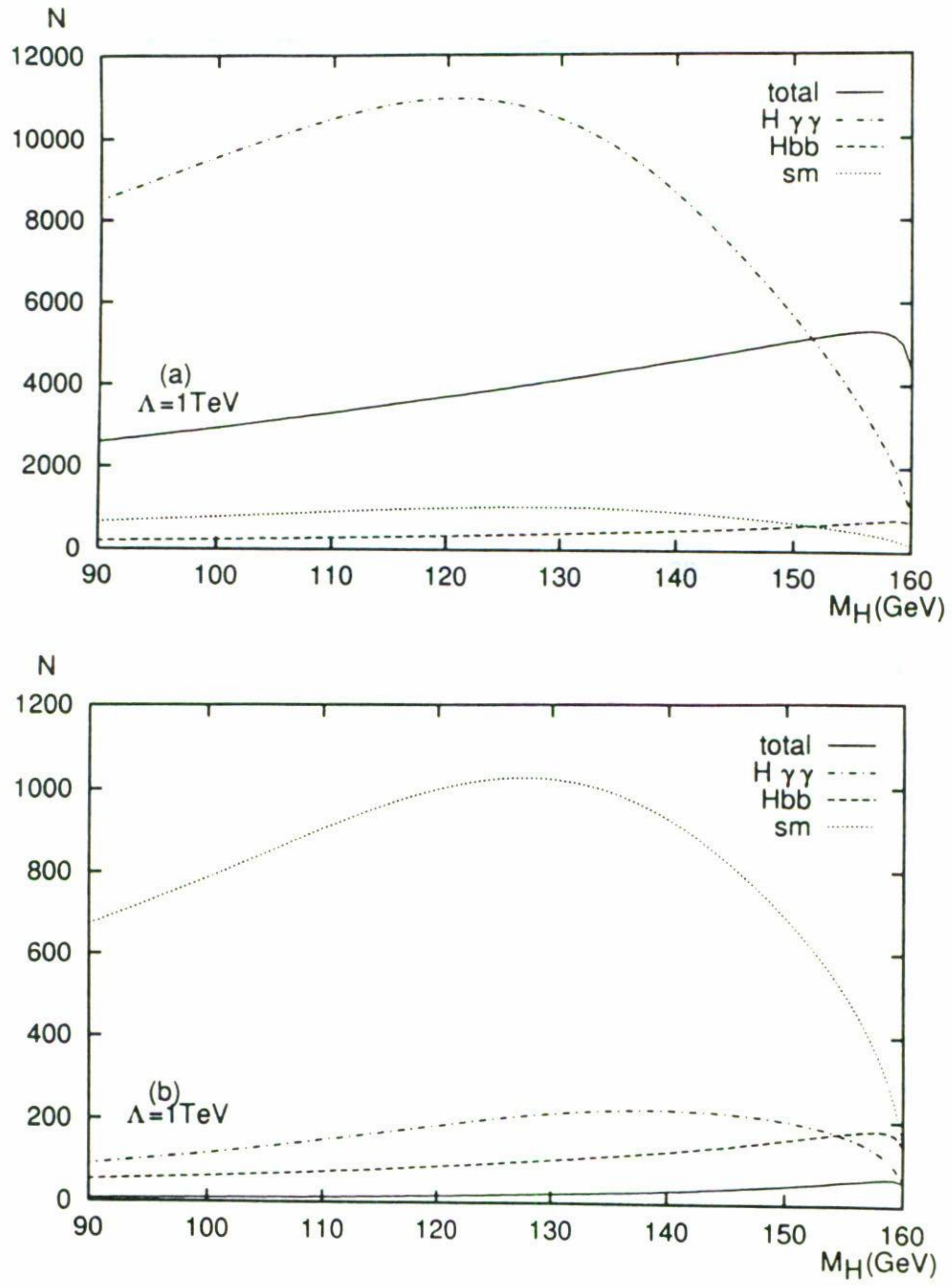

FIGURE 3. Number of events expected for the $\bar{b}b$ final state as a function of the Higgs boson mass and $\Lambda = 1$ TeV, for the scenarios with a maximum (a) and minimum (b) number of events. In each case $H^0\gamma\gamma$ (dot-dashed), $H^0\bar{b}b$ (dashed), SM (dotted) and total (solid line) contributions are displayed.

In Figs. 3–6 we display the number of expected events as a function of the Higgs boson mass for the final states $\bar{b}b$, WW, ZZ and $\bar{t}t$. In the case of the final state $\bar{b}b$ ($m_W < m_H < 2m_W$) we have also included the contributions of the WW^* and ZZ^* states

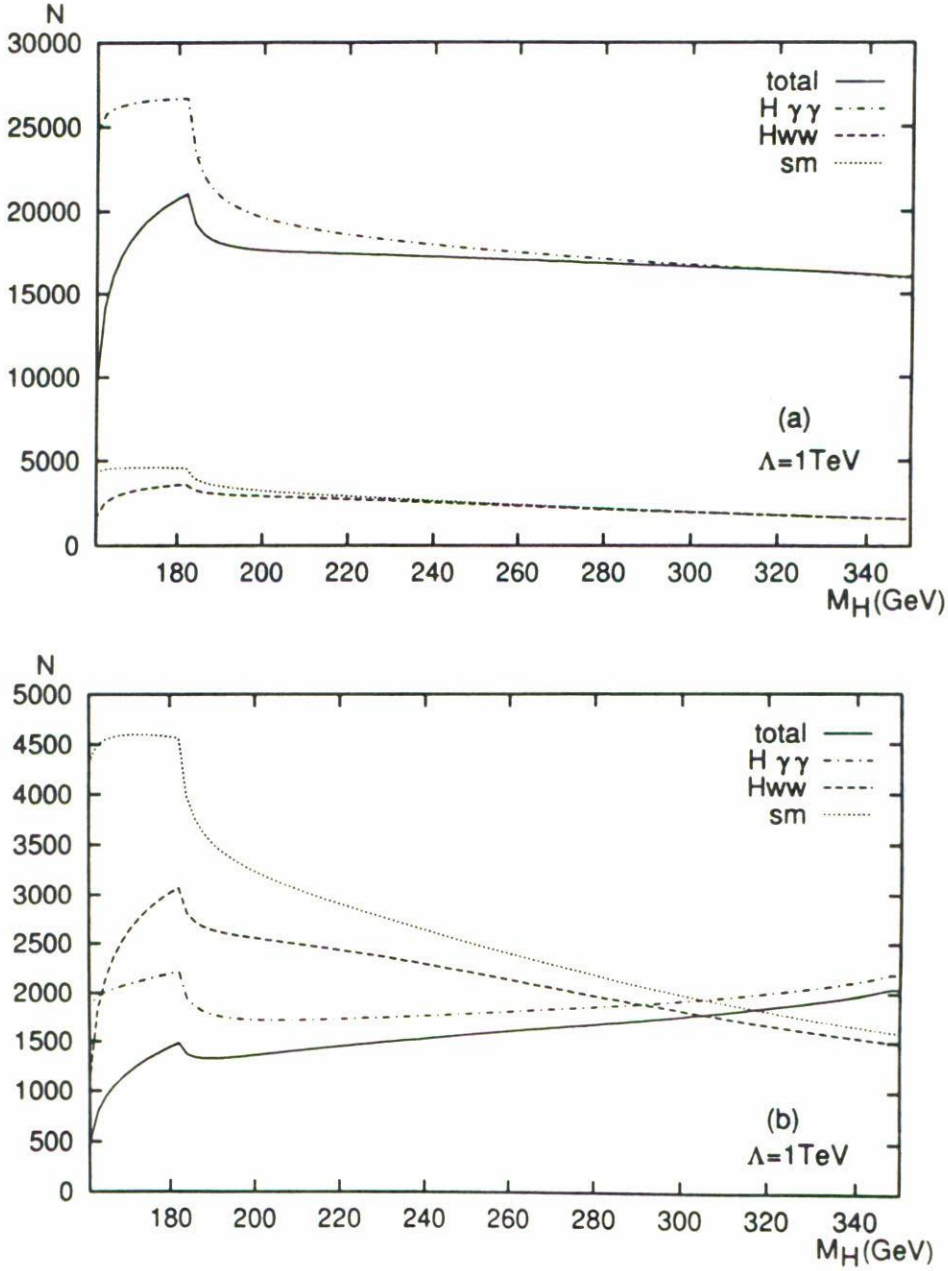

FIGURE 4. Number of events expected for the WW final state as a function of the Higgs boson mass and $\Lambda = 1$ TeV, for the scenarios with a maximum (a) and minimum (b) number of events. In each case, $H^0\gamma\gamma$ (dot-dashed), H^0WW (dashed), SM (dotted) and total (solid line) contributions are displayed.

on BR($H^0 \to \bar{b}b$). We have taken all the coefficients of the nonrenormalizable interactions as $|\alpha_i| = 1$ and the new physics scale as $\Lambda = 1$ TeV. We consider a center of mass energy of the e^+e^- collider equal to 500 GeV. The luminosity of the e^+e^- beam is taken equal

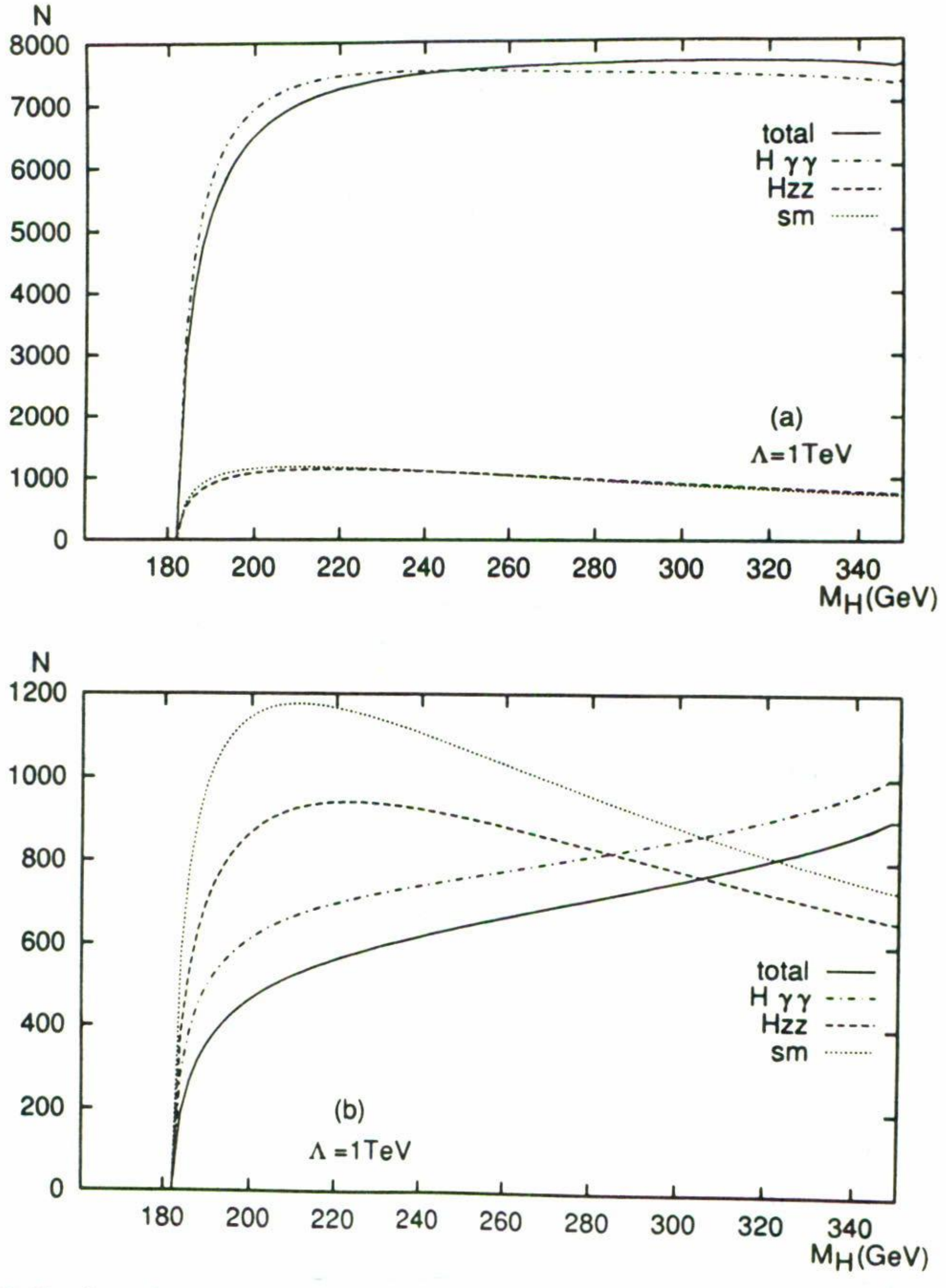

FIGURE 5. Number of events expected for the ZZ final state a a function of the Higgs boson mass and $\Lambda = 1$ TeV, for the scenarios with a maximum (a) and minimum (b) number of events. In each case, $H^0\gamma\gamma$ (dot-dashed), H^0ZZ (dashed), SM (dotted) and total (solid line) contributions are displayed.

to 20 fb^{-1}. Furthermore, we have considered the theoretical limit $\lambda\tilde{\lambda} = 1$. There are sixty-four different possibilities for the channels $\bar{b}b$, WW, and ZZ, and one-hundred and twenty-eight for the $\bar{t}t$ channel compatible with the above constraint. We present only the

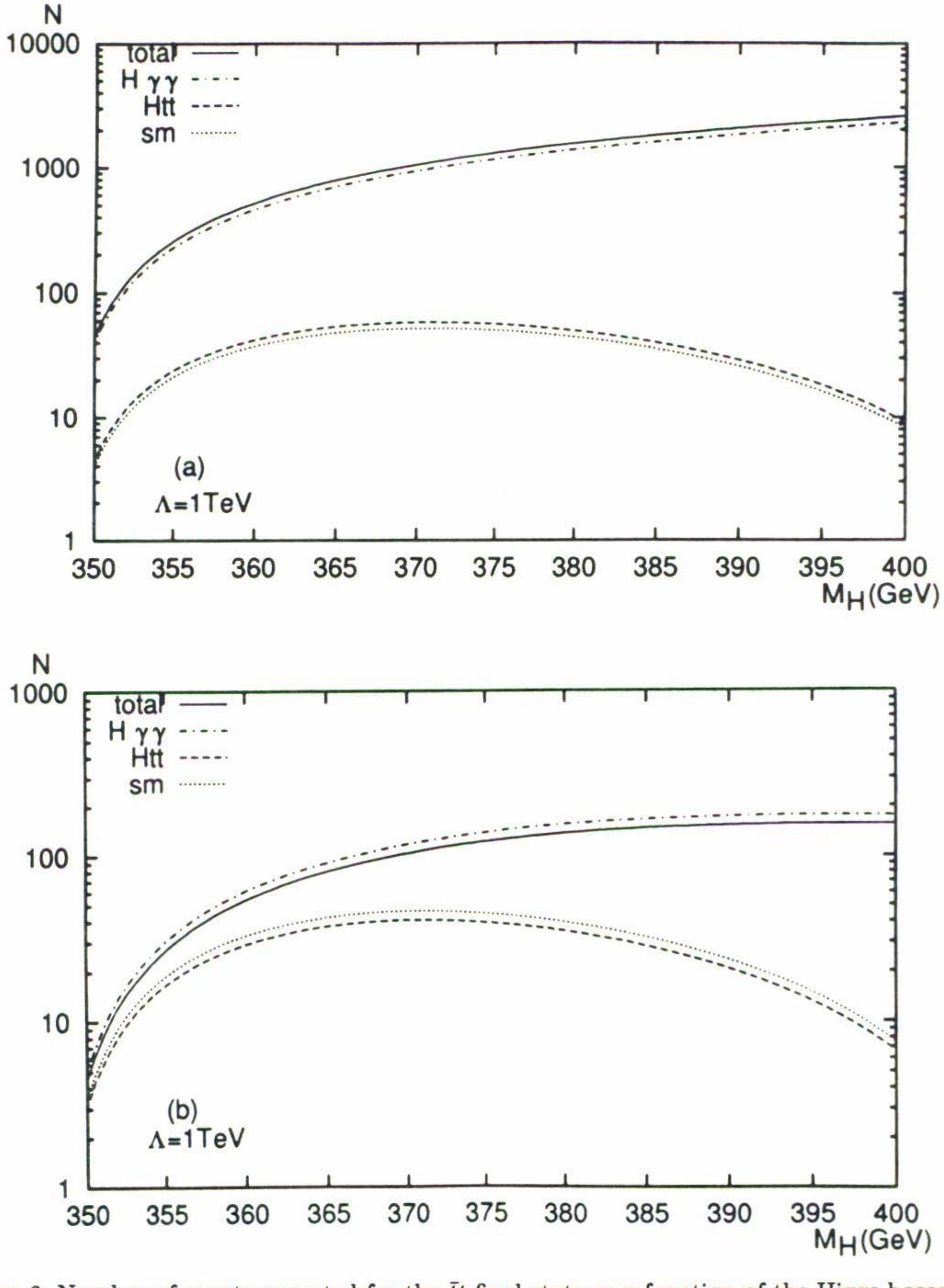

FIGURE 6. Number of events expected for the $\bar{t}t$ final state as a function of the Higgs boson mass and $\Lambda = 1$ TeV, for the scenarios with a maximum (a) and minimum (b) number of events. In each case, $H^0\gamma\gamma$ (dot-dashed), $H^0\bar{t}t$ (dashed), SM (dotted) and total (solid line) contributions are displayed. We have used a logarithmic scale in the vertical axis.

two extreme scenarios corresponding to a maximum and a minimum number of events for the total contribution, which includes effective interactions in all the vertices involved. Additionally, we have checked that for $\Lambda \geq 3$ TeV, the heavy physics effects decouples from the SM predictions. In Figs. 3-6, we display separately the SM contributions, the partial effective contributions due to new physics effects on $\Gamma(H^0 \to \gamma\gamma)$ and $\mathrm{BR}(H^0 \to X)$ ($X = \bar{b}b$, WW, ZZ and $\bar{t}t$), and the total contribution. In the scenario with a maximum number of events we can appreciate an enhancement of almost one order of magnitude, with respect to the SM predictions, for the final states $\bar{b}b$, WW and ZZ, while for the $\bar{t}t$ channel the enhancement may be up to two orders of magnitude. In the less favorable scenario the $\bar{b}b$ final state is suppressed by one order of magnitude with respect to the SM prediction, while in the $\bar{t}t$ final state subsists a remarkable enhancement with respect to the SM predictions [14].

In conclusion, the Higgs boson production through $\gamma\gamma$ fusion results in a highly sensitive mechanism to detect physics beyond the SM. Furthermore, we expect that the enhancement induced by new physics on the signal is not hidden by corresponding effects on the background [14].

4. ANOMALOUS EM COUPLINGS OF THE W BOSON AND QUARKS

We consider the following dimension-six operators which are $\mathrm{SU}(2)_L \times U(1)_Y$ gauge invariant and contribute to the flavor-changing neutral current (FCNC) decay $b \to s\gamma$ at tree- and one-loop level within the framework of effective Lagrangians [1]:

$$\begin{aligned} O_W &= \varepsilon_{ijk} W_\mu^{i\nu} W_\nu^{j\lambda} W_\lambda^{k\mu}, \quad & O_{WB} &= \phi^+ \tau^i \phi W_{\mu\nu}^i B^{\mu\nu}, \\ O_{uW}^{ab} &= \bar{Q}_L^a \sigma^{\mu\nu} W_{\mu\nu}^i \tau^i \tilde{\phi} u_{\mathrm{R}}^b, \quad & O_{dW}^{ab} &= \bar{Q}_L^a \sigma^{\mu\nu} W_{\mu\nu}^i \tau^i \phi d_{\mathrm{R}}^b, \\ O_{uB}^{ab} &= \bar{Q}_L^a \sigma^{\mu\nu} B_{\mu\nu} Y \tilde{\phi} u_{\mathrm{R}}^b, \quad & O_{dB}^{ab} &= \bar{Q}_L^a \sigma^{\mu\nu} B_{\mu\nu} Y \phi d_{\mathrm{R}}^b. \end{aligned} \tag{14}$$

After spontaneous symmetry breaking the fermionic operators generate effective vertices proportional to the anomalous magnetic moments of quarks and FCNC transitions like $b \to s\gamma$ and $t \to c\gamma$. Calculating the contribution of the effective vertices to the $b \to s\gamma$ transition in the Feynman-'t Hooft gauge with dimensional regularization will involve the one-loop diagrams shown in Fig. 7. A heavy dot denotes one of these vertices and there are additional graphs where the W bosons are replaced by the corresponding would-be Goldstone boson G_W.

In order to obtain the $b \to s\gamma$ branching fraction $\mathrm{Br}(b \to s\gamma)$, the inclusive quark-level process $b \to s\gamma$ is scaled to that of the semileptonic decay $b \to c\ell\nu$, including phase space and QCD corrections [15]. The calculation of the new contributions to $\Gamma(b \to s\gamma)$ due to the effective couplings involves only the dipole $b \to s$ transition operators denoted by O_7 [16]. The coefficient of this operator at the W scale can be expressed in the form $C_7(M_W) = C_7(M_W)_{\mathrm{SM}} + \Sigma\varepsilon_i \Delta C_7(M_W)_i$, where ε_i are the coefficients of each dimension-6 operator given in (14). Only the functions $\Delta C_7(M_W)_i$ for ε_W and ε_{WB} are finite and cutoff-independent. The respective functions for ε_{qW}^{ab} and ε_{qB}^{ab} involve a logarithmic dependence

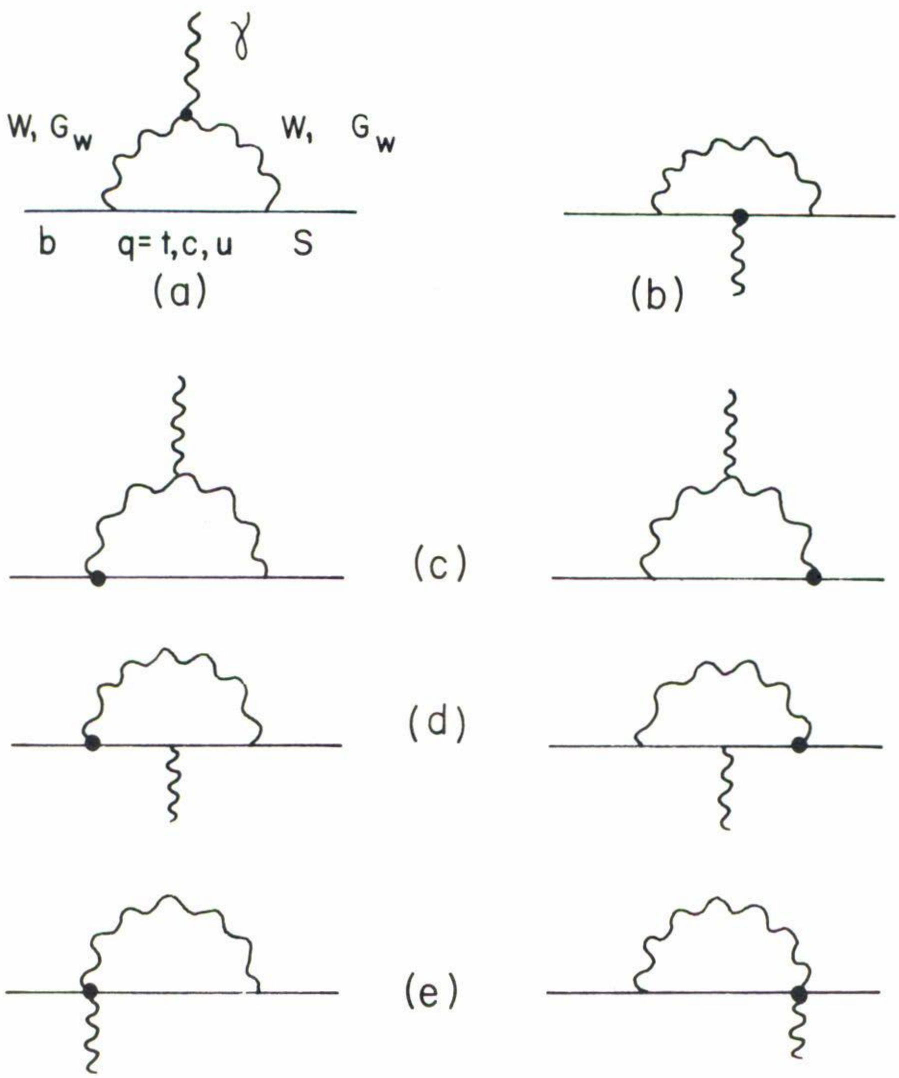

FIGURE 7. Feynman diagrams contributing to the $b \to s\gamma$ transition. The heavy dot denotes an effective vertex. There are additional graphs when the W lines are replaced by the corresponding would be Goldstone boson G_W.

on Λ. Figures 8 and 9 display $\mathrm{Br}(b \to s\gamma)$ as a function of the ε_W and ε_{WB} couplings for several top-quark masses assuming all other effective couplings are zero. In the SM contribution to $\Gamma(b \to s\gamma)$ we have included the next-to-leading log evolution equations for the coefficient of the $b \to s$ transition operators in the effective Hamiltonian [17], the gluon bremsstrahlung corrections [18], the $m_t \neq m_W$ corrections [19] and a running α_{QCD} evaluated at the b-quark mass scale [20].

The coefficients ε_W, ε_{WB}, ε_{qW}^{ab} and ε_{qB}^{ab} of the dimension–6 operators (14) are related to the $\delta\kappa_W$, $\delta\kappa_q$, λ_W parameters used to denote deviations of the magnetic dipole moments of the W boson and the quarks and the electric quadrupole moment of the W boson in

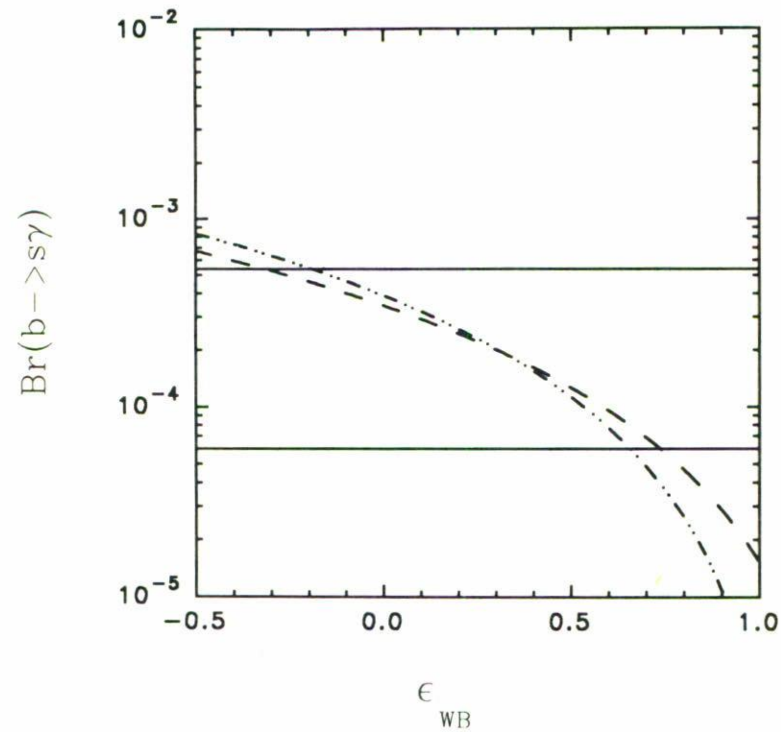

FIGURE 8. The branching franction $\mathrm{Br}(b \to s\gamma)$ as a function of ε_{WB} for two values of m_t, 160 GeV (dashed), and 200 GeV (dot-dashed). The bounds of the CLEO Collaborations are explicitly indicated.

the following form [21]:

$$
\begin{aligned}
\delta\kappa_W &= \kappa_W - 1 = \frac{g}{g'}\varepsilon_{WB}, \\
\lambda_W &= \tfrac{3}{2} g \varepsilon_W, \\
\delta\kappa_u^a &= -\sqrt{2}\frac{m_u}{m_W}\frac{g}{eQ_u}(c_W\varepsilon_{uB}^{aa} + s_W\varepsilon_{uW}^{aa}), \\
\delta\kappa_d^b &= -\sqrt{2}\frac{m_d}{m_W}\frac{g}{eQ_d}(c_W\varepsilon_{dB}^{bb} + s_W\varepsilon_{dW}^{bb}),
\end{aligned}
\tag{15}
$$

where $c_W = \cos\theta_W$, $s_W = \sin\theta_W$ and the SM limit is obtained with $\delta\kappa_W = \lambda_W = \delta\kappa_q = 0$.

The constraints on the ε_W and ε_{WB} anomalous couplings are obtained from the bounds placed by the CLEO Collaboration on the inclusive quark-level process $0.65 \times 10^{-4} < \mathrm{Br}(b \to s\gamma) < 5.4 \times 10^{-4}$ at 95% CL [22]. In each case, we assume that only one of these couplings is non-zero. We can appreciate that our bounds on ε_W and ε_{WB} as obtained from Figs. 8 and 9 do not improve analogous bounds obtained from low-energy precision data by de Rújula et al. [1]. Using Eqs. (15) our bounds can be put as $-0.41 < \delta\kappa_W < 1.22$ and $-2.31 < \lambda_W < 1.12$ for $m_t = 180$ GeV. These bounds agree with those obtained in the

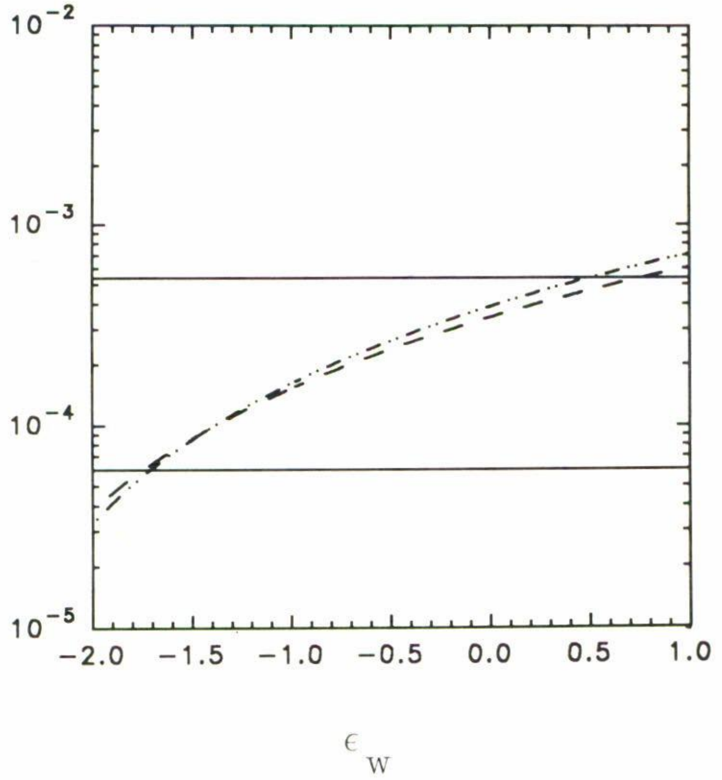

FIGURE 9. The branching fraction $\mathrm{Br}(b \to s\gamma)$ as a function of ε_W with the same conventions as in Fig. 8.

$U(1)_{EM}$–gauge invariant approach [23]. This is an expected result since both calculations differed only in the use of a specific gauge [24].

In Fig. 10 we display the $b \to s\gamma$ branching fraction as a function of $\delta\kappa_u$, $\delta\kappa_c$, $\delta\kappa_t$, for two different top-quark masses, assuming all other effective couplings are zero. Our bound for $\delta\kappa_t$ [24], the anomalous magnetic dipole moment of quark t, has improved by about one order of magnitude the $U(1)_{EM}$ gauge-invariant result [25]. In this case we expect some differences with respect to the $U(1)_{EM}$ gauge invariant case, essentially because the fermionic operators given in (14) induce contributions which are not present in the $U(1)_{EM}$ approach [25]. On the other hand, our bounds for the anomalous magnetic dipole moments of the lighter quarks u and c are weaker that those obtained "indirectly" from the Z decay widths reported at LEP-I [21]. The latter situation is a consequence of the quark mass dependence involved in the definition of $\delta\kappa_u$ and $\delta\kappa_c$ given in (15).

5. A MODEL INDEPENDENT BOUND ON $t \to c\gamma$

The recent detection of the semileptonic t-quark decays by the CDF and DO collaborations [26] puts as one of the most important tasks of future experiments to investigate carefully the production and decays of top quarks. In particular, it has been pointed

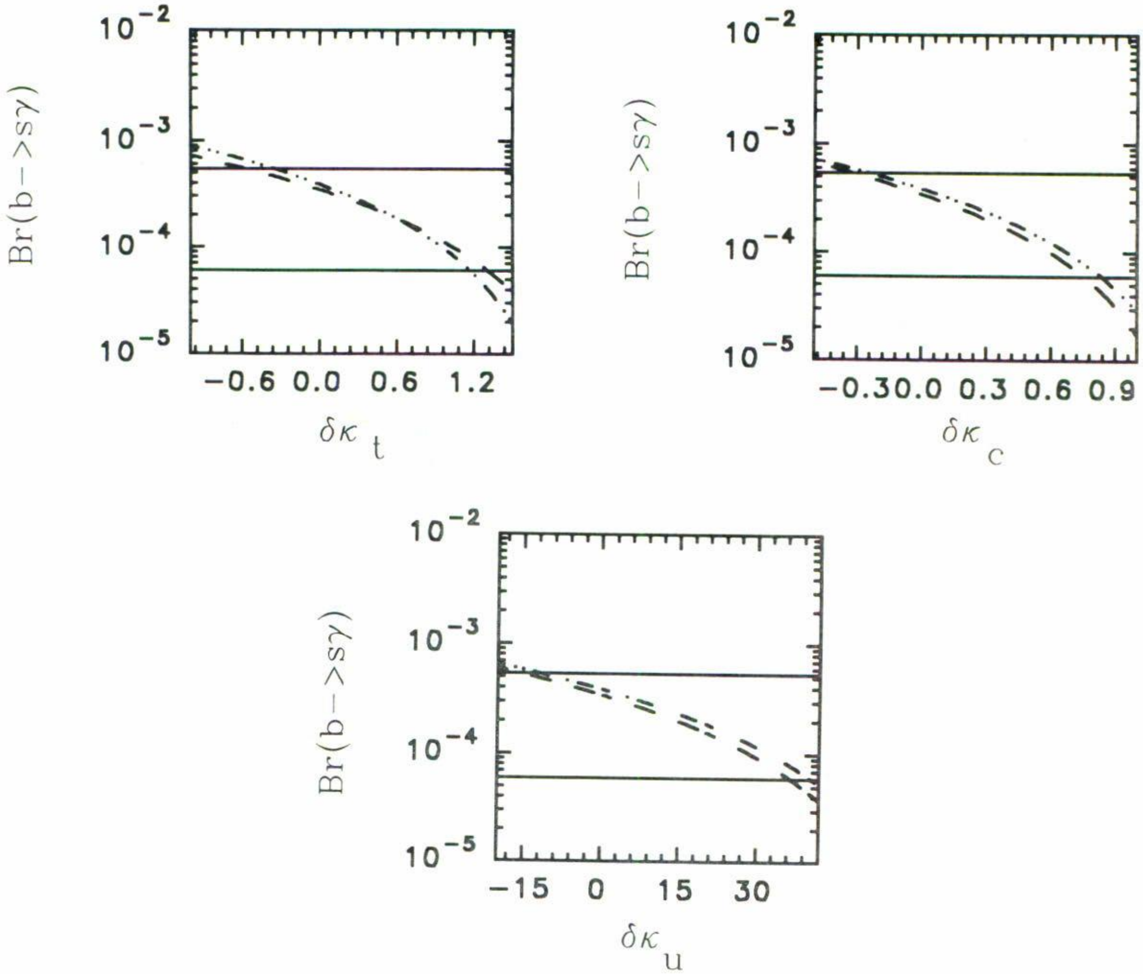

FIGURE 10. The branching fraction $\mathrm{Br}(b \to s\gamma)$ as a function of the anomalous magnetic moments of the quarks t, c, u. The conventions of Fig. 8. are used.

out [27] that the rare decay mode $t \to c\gamma$, although anticipated to be small, should provide a new window in the quest for physics beyond the SM. In the SM, the flavor-changing transition $t \to c\gamma$ has been found to be too small to be detectable due to a strong GIM suppression as well to the larger tree-level rate for $t \to Wb$ [27]. It has been found, however, that new contributions induced by charged Higgs bosons and charginos in the minimal supersymmetric standard model (MSSM) [28], or pseudo-Goldstone bosons in the one-generation technicolor model [29], may enhance the SM branching fraction by as much as 3–4 orders of magnitude.

In this section we examine the prospect of using the one-loop induced decay $b \to s\gamma$ as a laboratory to constrain the top quark rare decay $t \to c\gamma$. The possibility of constraining new physics effects by precision low-energy measurements in a model-independent manner has been stressed in the effective Lagrangian approach [1]. In this formalism, the weak radiative decays $b \to s\gamma$ and $t \to c\gamma$ are related through the Feynman diagrams shown in

Figs. 11(a) and 11(b). Here, the heavy dots denote effective $qqW(\gamma)$ couplings which arise from the dimension-six, $SU(2)_L \times U(1)$ gauge invariant operators given in (14).

Each of the effective operators given in (14) appears in the effective Lagrangian multiplied by undetermined coefficients ε^{ab}_{qW}, ε^{ab}_{qB}. The dependence on the Λ parameter is masked in these coefficients. The diagonal coefficients $\varepsilon^{a}_{qW(qB)}$ are related to the possible deviations of the magnetic dipole moments of quarks. On the other hand, the off-diagonal effective transitions $b \to s\gamma$ and $t \to c\gamma$ are proportional to the coefficients $\varepsilon^{23}_{dW(dB)}$ and $\varepsilon^{23}_{uW(uB)}$, respectively. Therefore, according to the effective Lagrangian approach, we can obtain a constraint on the rare decay transition $t \to c\gamma$ if we use a bound on the coefficient $\varepsilon^{23}_{dW(dB)}$ obtained from the CLEO result on $\mathrm{Br}(b \to s\gamma)$ and insert it in the Feynman diagrams shown in Figs. 11(b-e). Similarly, we can obtain another bound on the other coefficient $\varepsilon^{23}_{uW(uB)}$ by requiring that the corresponding contribution to diagrams 11(b)–(h) is consistent with the CLEO constraint through the dipole transition $b \to s$ coefficient $C_7(m_W) = C_7(m_W)_{SM} + \varepsilon^{23}_{uW(uB)}\Delta C_7(m_W)$ [16].

It is important to notice that diagrams 11(b–e) are not suppressed by the CKM matrix elements V_{tb} and V_{cs}. Furthermore, the calculation of these Feynman diagrams within the effective Lagrangian approach do not involve any other unknown parameter than the coefficients $\varepsilon^{ab}_{qW(qB)}$. As a consequence, we expect a rather transparent bound on $\mathrm{Br}(t \to c\gamma)$. Based on the recent bound for $b \to s\gamma$ obtained by the CLEO Collaboration [22], a remarkably high constraint is found for the FCNC top quark branching fraction, $\mathrm{Br}(t \to c\gamma) < 2 \times 10^{-5}$, which depends rather weakly on the top quark mass, $m_t \approx 180$ GeV [24].

On the other hand, using the constraint on the effective coefficient $\varepsilon^{23}_{dW(dB)}$ obtained directly from the $b \to s\gamma$ bound, $|\varepsilon^{23}_{dW(dB)}| < 1.69 \times 10^{-7}$, we obtain in this case the result $\mathrm{Br}(t \to c\gamma) < 4 \times 10^{-24}$, which is quite a small value. The effective Lagrangian approach leads to a direct interpretation of the above unequal bounds for $\mathrm{Br}(t \to c\gamma)$. While diagram 11(b) with the $tc\gamma$ insertion, together with the CLEO constraint $b \to s\gamma$, leads to a bound on the decay $t \to c\gamma$ through the effective tree-level coupling $t\gamma c$, diagram 1(b) leads to a bound on the same decay $t \to c\gamma$, but this time through an effective one-loop level contribution. The insertion of the effective $b\gamma s$ coupling in the latter case will result in a contribution equivalent to a two-loop level effect. This explains the high suppression obtained on the branching ratio calculated from diagram 11(b) when the quarks t and c are in the external legs.

In conclusion, the effective Lagrangian formalism and the CLEO constraint on $b \to s\gamma$ generate a model independent bound $\mathrm{Br}(t \to c\gamma) \leq 2 \times 10^{-5}$ for $m_t \approx 180$ GeV, which is consistent with $\mathrm{Br}(t \to c\gamma) \sim 10^{-12}$ obtained for the SM [27] and $\mathrm{Br} \sim 10^{-7}$–10^{-8} for SUSY and technicolor models [28, 29].

6. CONSTRAINTS ON $\nu_i \to \nu_j\gamma$

During the last few years, the electromagnetic properties of neutrinos have received increased interest [30]. In particular, it is known [31] that both Dirac and Majorana neutrinos can have off-diagonal (transition) magnetic moments κ_{ij} which become operative in the

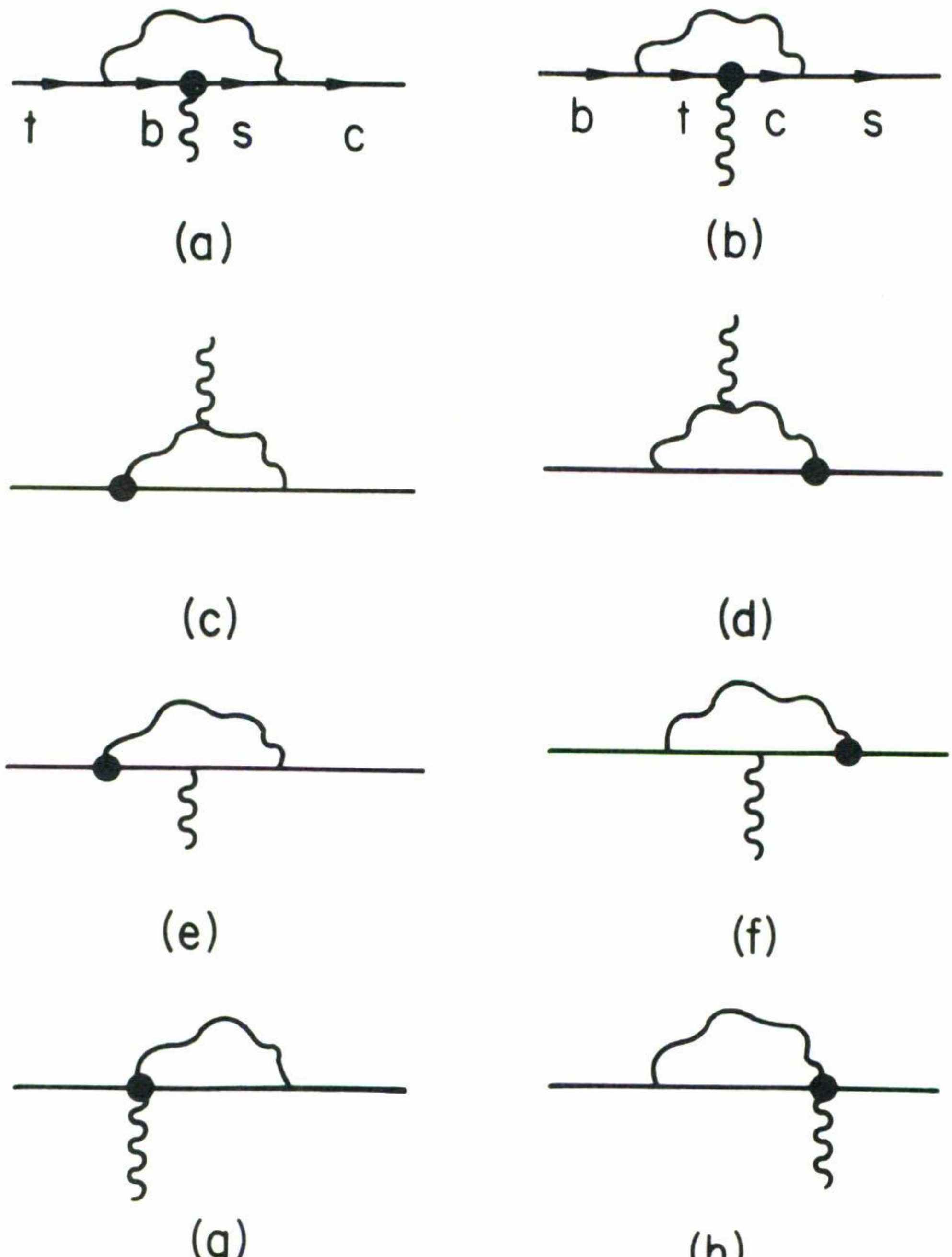

FIGURE 11. Feynman diagrams contributing to the $t \to c\gamma$ and $b \to s\gamma$ transitions. The heavy dot denotes an effective vertex. There are additional graphs when the W lines are replaced by the corresponding would be Goldstone boson G_W.

radiative decay $\nu_i \to \nu_j\gamma$. This decay mode is of particular interest in studying the transit of neutrinos from stellar interiors to the surface. Neutrino pair emission from white dwarfs and red giants sets the bound, for neutrino masses less than 10 KeV, $|\kappa_{e\mu}| \leq 10^{-11}$ in units of Bohr magnetons $(e/2m_e)$ [31]. Also, laboratory limits to $\nu_\mu e \to \bar{\nu}_e e$ and $\nu_e e \to \bar{\nu}_\mu e$ scattering via neutrino transition magnetic moments lead to the bound $|\kappa_{e\mu}| \leq 10^{-10}$ [30, 31].

We are interested in obtaining bounds on the lepton-number-nonconserving decays $\nu_\mu \to \nu_e\gamma$ and $\nu_\tau \to \nu_\mu\gamma$, $\nu_e\gamma$ from the known experimental constraints [32] on $\mu \to e\gamma$ and $\tau \to \mu\gamma, e\gamma$. For the sake of definiteness, we will consider Dirac neutrinos with finite rest mass in the framework of the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ SM. In the effective Lagrangian approach [1], dimension-six operators contribute to the latter radiative processes through the one-loop graphs (t'Hooft-Feynman gauge) shown in Fig. 12. The heavy dots indicate the dimension-six modifications to the SM low-energy Lagrangian. The anomalous vertices $W\ell_i\nu_j$, $W\ell_i\nu_j\gamma$ and $\nu_i\nu_j\gamma$ arise from the dimension-six operators

$$\begin{aligned} O_{\nu W}^{ab} &= \bar{\ell}_L^a\, \sigma^{\mu\nu}\, W_{\mu\nu}^i\, \tau^i \tilde{\phi}\, \nu_{\mathrm{R}}^b, \\ O_{\nu B}^{ab} &= \bar{\ell}_L^a\, \sigma^{\mu\nu}\, B_{\mu\nu}\, \tilde{\phi}\, \nu_{\mathrm{R}}^b, \end{aligned} \tag{16}$$

which preserve the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge symmetry of the SM.

The construction of the effective Lagrangian includes the dimension four terms corresponding to the SM contributions. There is one dimension five operator consisting of two lepton and two Higgs fields which violates lepton number conservation. However, such operator gives rise to a rather large Majorana mass for the neutrinos unless Λ, the energy scale representing the onset of new physical phenomena, is of order 10^{13} GeV, and hence uninteresting for the present study of the lepton-number nonconserving decays [33]. Therefore, we have to consider only dimension-six operators and all the effects from heavy excitations will be suppressed by at least two powers of Λ [1].

The transition magnetic moments, in units of Bohr magnetons, for neutrinos are given in terms of these coefficients ϵ^{ij} in the following form:

$$\kappa_{ij} = \sqrt{2}\frac{m_e}{m_W}\frac{g}{e}(\sin\theta_{\mathrm{W}}\epsilon_{\nu W}^{ij} + \cos\theta_{\mathrm{W}}\epsilon_{\nu B}^{ij}). \tag{17}$$

Therefore, a constraint on the coefficients ϵ^{ij} obtained from the experimental bounds on $\Gamma(\ell_i \to \ell_j\gamma)$ will be translated into a constraint on the transition magnetic moments κ_{ij}. Calculating the contribution of the effective operators (16) to these flavor changing transitions in the t'Hooft-Feynman gauge with dimensional regularization will involve the one-loop diagrams shown in Fig. 7.

The decay width for the process $\mu \to e\gamma$ can be expressed as

$$\Gamma(\mu \to e\gamma) = \frac{1}{8\pi}\left(\frac{m_\mu^2 - m_e^2}{m_\mu}\right)^3 (m_e^2 + m_\mu^2)\left|\epsilon_{\nu W}^{12}\, F_{2\mathrm{R}}^W(m_\nu) + \epsilon_{\nu B}^{12}\, F_{2\mathrm{R}}^B(m_\nu)\right|^2, \tag{18}$$

where the functions $F_{2\mathrm{R}}^{B,W}$ are defined in terms of the transition amplitude

$$M = \bar{u}_e(p)\, i\sigma_{\mu\nu}\, q^\nu\, \epsilon^\nu(m_\mu P_{\mathrm{R}} + m_e P_{\mathrm{L}})u_\mu(p)(\epsilon_{\nu W}^{12}\, F_{2\mathrm{R}}^W + \epsilon_{\nu B}^{12}\, F_{2\mathrm{R}}^B), \tag{19}$$

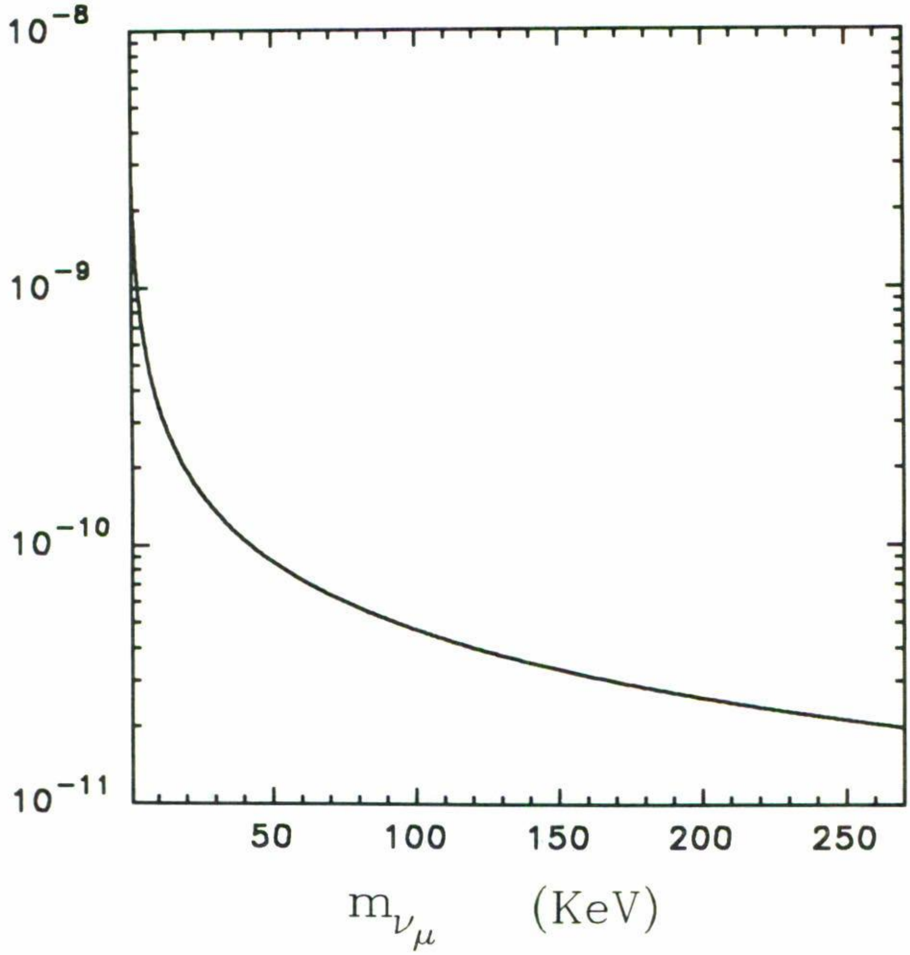

FIGURE 12. The bound for the transition magnetic moment $\kappa_{\mu e}$ as a function of the muon neutrino mass.

with $q = p - p'$, $P_{\mathrm{R,L}} = (1 \pm \gamma_5)/2$, and ϵ^ν being the photon polarization. In the approximation $m_{\nu_e} \ll m_{\nu_\mu}$ we get [34]

$$F_{2\mathrm{R}}^W \simeq 4G_{\mathrm{F}}\, e \frac{m_W m_{\nu_\mu}}{m_\mu^2} \left[\frac{3}{2} + \ln \frac{m_{\nu_\mu}^2}{m_\mu^2} + 3m_W^2 C_0(p, p-q, m_{\nu_\mu}, m_W, m_W)\right], \qquad (20)$$

$$F_{2\mathrm{R}}^B \simeq 4G_{\mathrm{F}}\, e \cot\theta_{\mathrm{W}} \frac{m_W m_{\nu_\mu}}{m_\mu^2} \ln \frac{m_{\nu_\mu}^2}{m_\mu^2}, \qquad (21)$$

where C_0 is the scalar three-point integral [35] for the one-loop diagram shown in Fig. 7. Also, we have neglected the electron mass whenever it is appropriate.

From our result given in (18) and the experimental bound [32] $\Gamma(\mu \to e\gamma) \leq 14.68 \times 10^{-30}$ GeV, we get a constraint on each coefficient ϵ^{12}. We consider one operator at a time in order to avoid unnatural cancellations of the effects produced by different operators.

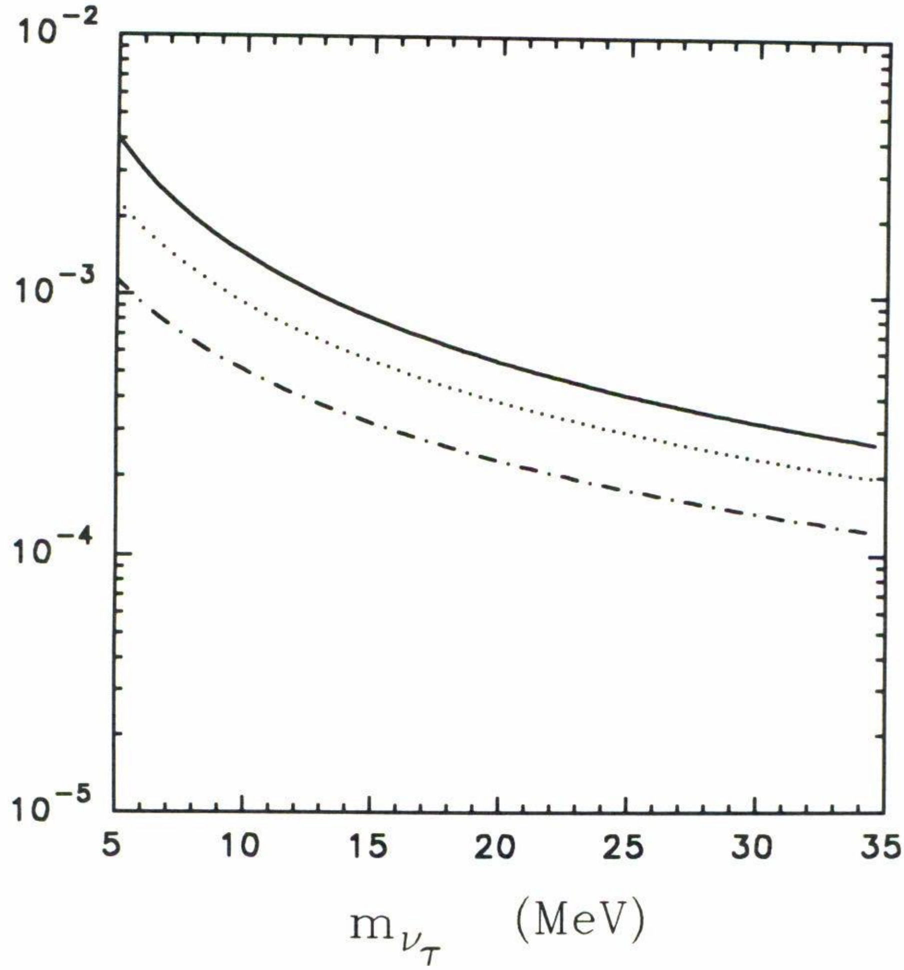

FIGURE 13. The bound for the transition magnetic moment $\kappa_{\tau\mu}$ as a function of the tau neutrino mass for three values of the ν_μ mass: 0.27 MeV (continuous line), 0.1 MeV (dot), and 0.01 MeV (dot-dash).

We then use (17) to obtain an upper bound for the transition magnetic moment. In Fig. 12 we plot this bound ($\kappa_{\mu e}$) as a function of the muon neutrino mass. In Figs. 13 and 14 we present the respective upper bounds for $\kappa_{\tau e}$ and $\kappa_{\tau\mu}$ obtained in a similar way. For the latter case we have used the full expressions for the parametric functions $F_{2R}^{B_\tau, W_\tau}$ in terms of the Veltman-Passarino functions B_0 and C_0 [35]. We have also used the experimental bounds [32] $\Gamma(\tau \to e\gamma) \leq 4.316 \times 10^{-16}$ GeV, $\Gamma(\tau \to \mu\gamma) \leq 11.87 \times 10^{-16}$ GeV with $m_{\nu_\mu} \leq 270$ KeV and $m_{\nu_\tau} \leq 35$ MeV. From our results shown in Figs. 12–14 it is easy to include more recent bounds for the neutrino masses [36].

The predicted rates for the decay width $\Gamma(\nu_i \to \nu_j \gamma)$ within the SM with massive neutrinos are rather small [37]. Models with charged Higgs bosons or a fourth generation of leptons may enhance these decay rates. In these cases it is possible to reach values for the transition magnetic moment of order $\kappa_{\mu e} \approx 10^{-12}$ or even higher [38]. However, these

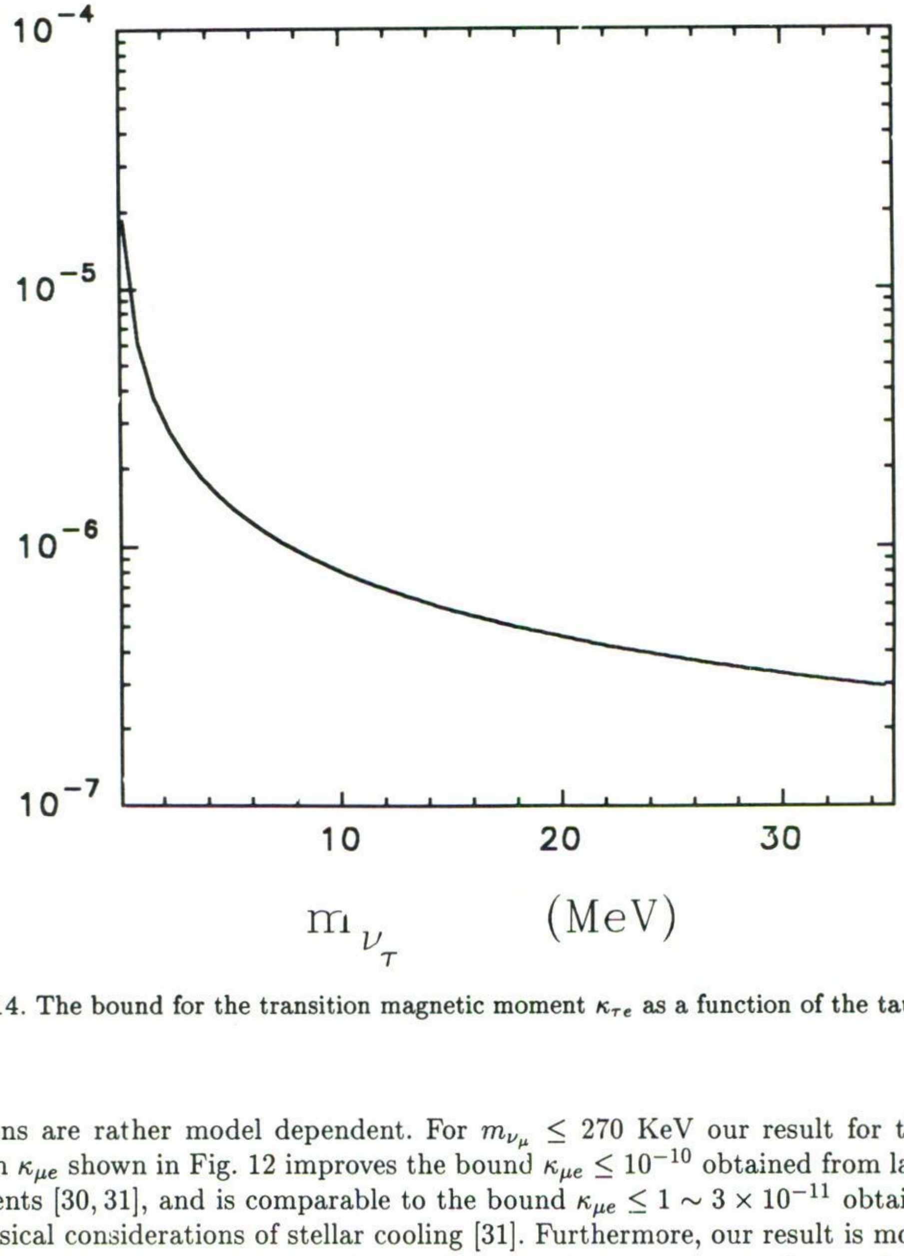

FIGURE 14. The bound for the transition magnetic moment $\kappa_{\tau e}$ as a function of the tau neutrino mass.

predictions are rather model dependent. For $m_{\nu_\mu} \leq 270$ KeV our result for the upper bound on $\kappa_{\mu e}$ shown in Fig. 12 improves the bound $\kappa_{\mu e} \leq 10^{-10}$ obtained from laboratory experiments [30, 31], and is comparable to the bound $\kappa_{\mu e} \leq 1 \sim 3 \times 10^{-11}$ obtained from astrophysical considerations of stellar cooling [31]. Furthermore, our result is model independent and it rests only on general considerations involved in the effective Lagrangian approach [1]. It is also important to mention that the upper bounds presented in Figs. 12–14 for the neutrino transition magnetic moment have a rather smooth dependence on the scale parameter Λ associated with the onset of new physical phenomena. This situation is similar to the one encountered for the constraints imposed by the $b \to s\gamma$ decay on the anomalous electromagnetic couplings of gauge bosons and quarks. A new upper bound for the transition magnetic moment of Majorana μ–neutrinos has been reported recently [39]. Eventhough the reported bound $\kappa_{\mu e} < 10^{-14}$ is rather restrictive, it rests again on astrophysical considerations. The analysis presented in this section can be applied to Majorana

μ–neutrinos within the context of left-right symmetric models. However, it was not possible to improve in this case the previous bounds for the neutrino transition magnetic moments [40].

We would like to point out that eventhough our results for the upper bounds on $\kappa_{\tau\mu}$ and $\kappa_{\tau e}$ displayed in Figs. 13 and 14 are weaker than the bound for $\kappa_{\mu e}$, as far as we know this is the first time that the off-diagonal magnetic moments of the τ neutrino are constrained in a model-independent scheme. Finally, these constraints can be translated into lower bounds for the respective lifetimes if use is made of (18): $\tau(\nu_\mu \rightarrow \nu_e\gamma) \geq 1.1 \times 10^4$ sec $\tau(\nu_\tau \rightarrow \nu_e\gamma) \geq 2.4 \times 10^{-11}$ sec and $\tau(\nu_\tau \rightarrow \nu_\mu\gamma) \geq 2.8 \times 10^{-17}$ sec [34].

7. Summary

We have employed the method of effective Lagrangians in evaluating the possible deviation of SM predictions on the weak radiative processes $H \rightarrow \gamma\gamma$, $b \rightarrow s\gamma$, $\mu \rightarrow e\gamma$ and on the production of Higgs bosons on $\gamma\gamma$ colliders. These calculations illustrate the fact that linear effective Lagrangians can be used in loop calculations [41].

We found important deviations stemming from the anomalous couplings when the scale of new physics is moderately large (a few TeV). In the $H \rightarrow \gamma\gamma$ decay and the production of Higgs bosons in $\gamma\gamma$ colliders we have found that the systematic bookkeeping of the α_i anomalous coefficients leads us to expect that an enhancement of these processes can possibly arise from the contribution of three-level-generated dimension-eight bosonic operators. However, in Ref. [11] it was obtained that the anomalous contributions to the decay widths $H^0 \rightarrow \gamma Z$ and $Z \rightarrow \gamma H$ are marginally important with respect to the $H \rightarrow \gamma\gamma$ decay mode.

On the other hand, the virtual effects of high-energy physics are usually broadly distributed. The search for them should concentrate in areas where their effects may be large. In particular, we have stressed the interesting observation that one can set model independent limits on the flavor changing transitions $t \rightarrow c\gamma$ and $\nu_i \rightarrow \nu_j\gamma$ from the CLEO measurement on $b \rightarrow s\gamma$ and the known experimental bound on $\mu \rightarrow e\gamma$. The effective Lagrangian approach relates the $tc\gamma$ and $\nu_i\nu_j\gamma$ effective couplings with the diagrams shown in Fig. 11 which gives a potential new physics contribution to $b \rightarrow s\gamma$ and $\mu \rightarrow e\gamma$ decays. The suppression of these contributions by an additional power of the weak coupling is partly compensated by the lack of the respective small CKM matrix elements. Therefore, an interesting limit on $t \rightarrow c\gamma$ and $\nu_j \rightarrow \nu_j\gamma$ are derived which compares favorably with the SM or SUSY expectations, or the scattering and astrophysical bounds in the case of the neutrino magnetic transition.

The conclusions of this paper are, by necessity, speculative. Still, based on the calculations presented it is clear that study of the processes we have dealt with may open a window into new physics.

Acknowledgments

We wish to thank J.L. Díaz Cruz, J.C. D'Olivo, J.G. Hirsh and J.W.F. Valle for discussions. This research was supported by CONACyT (México) and COLCIENCIAS (Colombia).

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