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Single top production at LEP II

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We study the production of a single top quark through the processes $e^+\gamma \to \nu t\bar{b}$, $e^+e^- \to W^*\gamma^* \to t\bar{b} + e\nu$, and $e^+e^- \to W^*W \to t\bar{b} + e\nu$ within the minimal Standard Model. Numerical estimates of the signal cross-sections at LEP II energies are compared with pair production below threshold, $e^+e^- \to t\bar{t}^* \to tW^*\bar{b} \to t\bar{b} + e\nu$. We show that if $m_{\rm top} > \sqrt{s}/2$, the WW^* and $W^*\gamma^*$ channels are dominant, but these processes give an interesting event rate only if the expected luminosity could be ameliorated relative to present estimates.

The top quark is the only fermion of the standard model still to be discovered. The present experimental lower bound on the top mass from CDF [1] is $m_{\rm top} \gtrsim 108$ GeV$/c^2$. In this letter we explore the possibility of producing the top quark in an e^+e^- machine such as LEP II (with energies in the range 200–250 GeV) in the event that $m_{\rm top} > \sqrt{s}/2$. In this case the real production process $e^+e^- \to t\bar{t}$ is forbidden by phase space considerations and the production of a $t\bar{t}$ pair must proceed with one or both top quarks off the mass shell. One expects the cross section for the process

$$e^+e^- \to t\bar{t}^* \to tW^*\bar{b} \to t\bar{b} + e\nu \tag{1}$$

to drop dramatically for $m_{\rm top} > \sqrt{s}/2$ [2]. In principle, a larger contribution can come through the process (W^*W channel)

$$e^+e^- \to W^*W \to t\bar{b} + e\nu\,, \tag{2}$$

where the e^+e^- incoming pair annihilates into one real and one virtual W boson. The real W is then observed to decay into an electron and a neutrino, while the other is produced off-the-mass-shell and converts itself into a real $t\bar{b}$ pair.

Finally we must also consider the process ($W^*\gamma^*$ channel)

$$e^+e^- \to W^*\gamma^* \to t\bar{b} + e\nu\,, \tag{3}$$

which bears some similarities to W–gluon fusion production of heavy quarks already studied [3] at hadron colliders. This process has been studied [4] for *massless quarks* because it is a potential background for the detection of an intermediate Higgs boson via its decay into $q\bar{q}$ pairs.

We report here our calculation of the cross section for the above processes at presently planned LEP II energies. For $m_{\rm top} < \sqrt{s}/2$, the tt^* channel obviously dominates and gives a substantial number of events. If on the other hand $m_{\rm top} > \sqrt{s}/2$, the WW^* channel is larger until it is surpassed by the $W^*\gamma^*$ channel. In all the three cases, however, we find that, with an integrated luminosity of $\mathcal{L} = 10^5$ nb^{-1}/yr, the event rate is still very low. Much

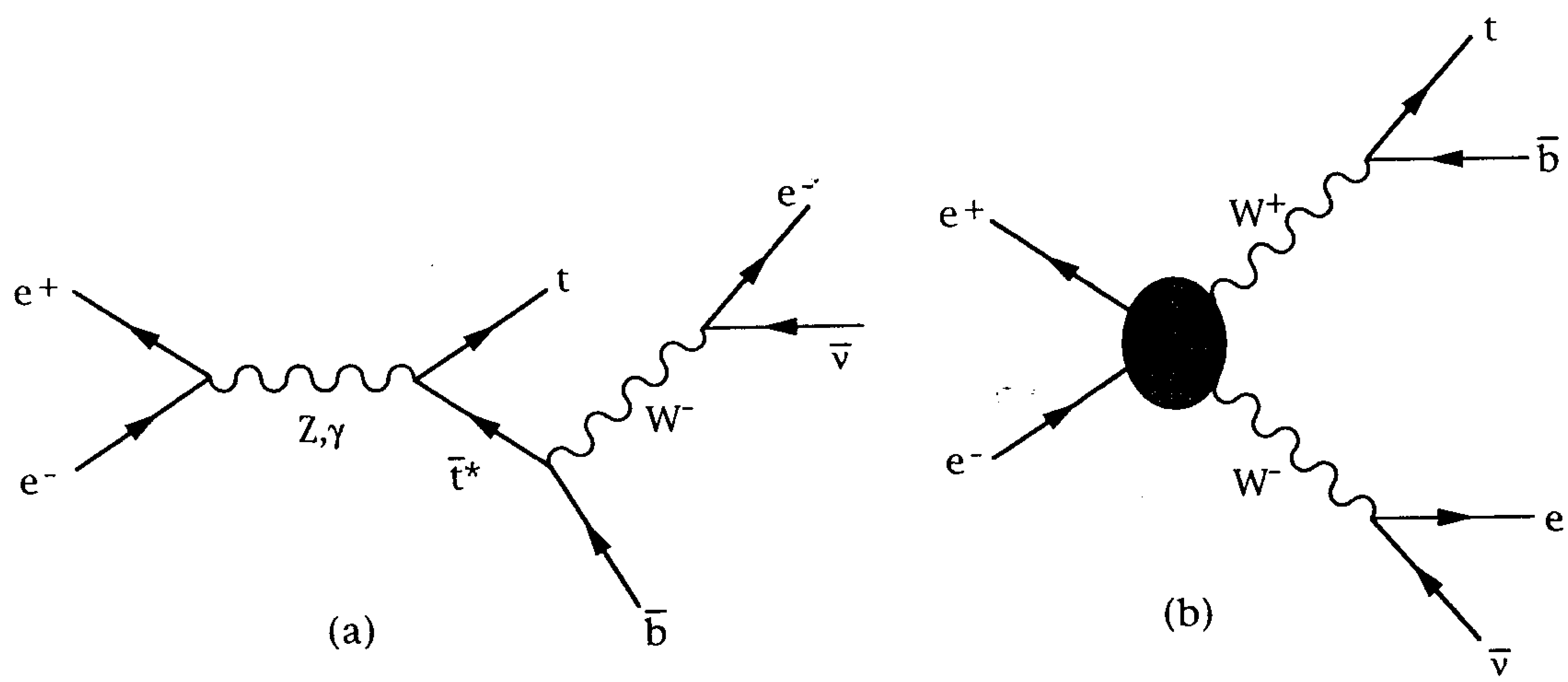

Fig. 1. (a) Production of a tt^* pair through the γ and Z in e^+e^- annihilation. (b) Production of the top quark through the decay of an off-the-mass-shell W.

higher luminosities are required before the WW^* and $W^*\gamma^*$ channels can contribute significantly to the detection of the top quark.

First we discuss the process of eq. (1) whose lowest order Feynman diagram is shown in fig. 1a. The amplitude is readily written down:

$$\mathcal{M} = \sum_{V=\gamma,Z} g_V\,\overline{v}(k_2)\gamma^\mu(a_V + b_V\gamma_5)u(k_1)$$
$$\times \frac{1}{s - M_V^2 + \mathrm{i}M_V\Gamma_V}\, g_V\,\overline{u}(p_1)\gamma_\mu(A_V + B_V\gamma_5)(-\not{p} + m_t)\gamma_\nu \frac{1-\gamma_5}{2} v(p_3)$$
$$\times \frac{g}{\sqrt{2}} V_{t,b} \frac{1}{Q^2 - m_t^2 + \mathrm{i}m_t\Gamma_{t^*}} \frac{g}{\sqrt{2}} \frac{1}{k^2 - M_W^2 + \mathrm{i}M_W\Gamma_W}\,\overline{u}(q)\gamma^\nu \frac{1-\gamma_5}{2} v(q'), \tag{4}$$

with the usual couplings $a_\gamma = -1$, $b_\gamma = 0$, $a_Z = -\frac{1}{4} + \sin^2\theta_w$, $b_Z = +\frac{1}{4}$, $A_\gamma = \frac{2}{3}$, $B_\gamma = 0$, $A_Z = \frac{1}{4} - \frac{2}{3}\sin^2\theta_w$, $B_Z = -\frac{1}{4}$, $g_\gamma = e$, $g_Z = g/\cos\theta_w = e/\cos\theta_w \sin\theta_w$ and $Q^2 = p^2$, $k^2 = (p-p_3)^2$ respectively the squared invariant masses of the virtual top quark and W-boson. A long, but straightforward calculation gives the differential cross section as a function of the invariant mass of the virtual top:

$$\left(\frac{\mathrm{d}\sigma}{\mathrm{d}Q^2}\right)^{tt^*} = \widetilde{\sigma}(s;m_t^2,Q^2)\,\frac{Q}{\pi}\,\frac{\Gamma(Q)}{(Q^2-m_t^2)^2 + (m_t\Gamma_{t^*})^2}. \tag{5}$$

The function $\widetilde{\sigma}$ may be regarded as the cross section for producing a $t(m_t^2)\,\overline{t}(Q^2)$ pair, and is given by

$$\widetilde{\sigma}(s;m_t^2,Q^2) = \sqrt{s}\sum_{V,V'=\gamma,Z}(a_Va_{V'} + b_Vb_{V'})\,\frac{g_Vg_{V'}S_{VV'}(s;m_t^2,Q^2)}{D_V(s)\,[D_{V'}(s)]^*}, \tag{6}$$

where the propagator for the gauge boson $V = \gamma, Z$ is $D_V(s) = s - M_V^2 + \mathrm{i}M_V\Gamma_V$ and $S_{VV'}$ is defined by

$$S_{VV'} = 3\,\frac{g_Vg_{V'}}{12\pi\sqrt{s}^3}\Big[3m_t^2(A_VA_{V'} - B_VB_{V'}) + \frac{1}{Q^2}\left(2s - Q^2 - m_t^2 - \frac{(Q^2-m_t^2)^2}{s}\right)$$
$$\times\left[\tfrac{1}{4}(Q^2+m_t^2)(A_VA_{V'} + B_VB_{V'}) - \tfrac{1}{4}(Q^2-m_t^2)(A_VB_{V'} + B_VA_{V'})\right]\Big]\lambda^{1/2}(s,m_t^2,Q^2). \tag{7}$$

$\Gamma(Q)$ is the partial width of the decay $t \to Wb \to e\nu b$ for a heavy top of mass Q which has already been discussed elsewhere [5]:

$$\Gamma(Q) = \Gamma(t^* \to Wb \to e\nu b)$$

$$= \frac{M_W^4 G_F^2 |V_{t,b}|^2}{96\pi^3 Q^3} \int_0^{(Q-m_b)^2} dk^2 \, \frac{\lambda^{1/2}(Q^2,k^2,m_b^2)}{(k^2-M_W^2)^2+(M_W\Gamma_W)^2} [(Q^2-m_b^2)^2 + k^2(Q^2+m_b^2) - 2k^4] . \quad (8)$$

The total width of a top quark of mass m_t is obtained summing over all the decay channels and, in the massless limit, is given by $9\Gamma(m_t)$. In the amplitude of eq. (4), Γ_{t^*} is assumed to be $9\Gamma(Q)$. It is easy to check that for $Q^2 = m_t^2$ the function $\widetilde{\sigma}$ reduces to the cross section [6] for the production of a real $t\bar{t}$ pair, i.e. $\lim_{Q^2 \to m_t^2} \widetilde{\sigma}(s; Q^2, m_t^2) = \sigma(e^+e^- \to t\bar{t})$.

Let us now turn our attention to the process of eq. (2) shown in fig. 1b. At the lowest order in perturbation theory, the amplitude is given by the three Feynman diagrams relative to $e^+e^- \to W^+W^-$ (see ref. [7]).

The lowest-order amplitude for the process $e^+e^- \to (W^+)^*(W^-)^* \to t\bar{b} + e^-\bar{\nu}$ is

$$\mathcal{M} = \bar{v}(k_2) T_{\mu\nu} u(k_1) \frac{\eta^{\mu\sigma} - p_2{}^\mu p_2{}^\sigma/M_W{}^2}{p_2^2 - M_W^2 + iM_W\Gamma_W} \frac{g}{\sqrt{2}} V_{tb} J_\sigma(2) \frac{\eta^{\nu\tau} - p_1{}^\nu p_1{}^\tau/M_W{}^2}{p_1^2 - M_W^2 + iM_W\Gamma_W} \frac{g}{\sqrt{2}} J_\tau(1) . \quad (9)$$

$T_{\mu\nu}$ is related to the amplitude for $e^+e^- \to W^+W^-$, which has been studied in detail by Brown and Mikaelian [7] and for the sake of comparison, in the following we shall use their notation. In eq. (9) $J(1)$ and $J(2)$, are the charged currents of the fermions through which the virtual W gauge bosons decay:

$$J_\tau(1) = \bar{u}(q_1)\gamma_\tau \frac{1-\gamma_5}{2} v(q_2) , \quad J_\sigma(2) = \bar{u}(q_3)\gamma_\sigma \frac{1-\gamma_5}{2} v(q_4) .$$

A simplification arises from the assumption that one of the W's (W^- for instance) decays into a light lepton pair ($e^-\bar{\nu}$). In this case, neglecting the electron mass, $q_1^2 = q_2^2 = 0$, so that $p_1 \cdot J(1) = 0$, the amplitude becomes

$$\mathcal{M} = \bar{v}(k_2) T_{\mu\nu} u(k_1) (g/\sqrt{2}) V_{tb} \frac{\eta^{\mu\sigma} - p_2{}^\mu p_2{}^\sigma/M_W{}^2}{p_2^2 - M_W^2 + iM_W\Gamma_W} J_\sigma(2) \frac{(g/\sqrt{2})}{p_1^2 - M_W^2 + iM_W\Gamma_W} J^\nu(1) . \quad (10)$$

The calculation of the cross section is carried out factorizing the Lorentz invariant four-particle phase space, and integrating over the momenta of the final state leptons q_1, q_2, q_3, q_4 to obtain ($Q_1^2 = p_1^2$ and $Q_2^2 = p_2^2$)

$$\left(\frac{d\sigma}{dQ_1^2 dQ_2^2}\right)^{WW^*} = \frac{1}{2s} \int d^2PS(K; p_1, p_2) \cdot \tfrac{1}{4} \mathrm{Tr}(\not{k}_2 T_{\mu\nu} \not{k}_1 \gamma_0 T^\dagger_{\mu'\nu'} \gamma_0) \frac{1}{2\pi} \frac{3(g/\sqrt{2})^2 |V_{tb}|^2}{|D_W(Q_2^2)|^2} \frac{1}{2\pi} \frac{(g/\sqrt{2})^2}{|D_W(Q_1^2)|^2}$$
$$\times \widetilde{T}^{\nu\nu'}(p_1; 0, 0)(\eta^{\mu\sigma} - p_2{}^\mu p_2{}^\sigma/M_W{}^2)(\eta^{\mu'\sigma'} - p_2{}^{\mu'} p_2{}^{\sigma'}/M_W{}^2) \widetilde{T}_{\sigma\sigma'}(p_2; m_t^2, m_b^2) , \quad (11)$$

with

$$\widetilde{T}_{\mu\mu'}(p; q^2, q'^2) = \int d^2PS(p; q, q') \, 2(q_\mu q'_{\mu'} + q_{\mu'} q'_\mu - q \cdot q' \eta_{\mu\mu'} - iq^\alpha q'^\beta \epsilon_{\alpha\beta\mu\mu'}) . \quad (12)$$

The result of the calculation of the phase space integration of eq. (12) is given by the following expression:

$$\widetilde{T}_{\mu\mu'}(p, q^2, q'^2) = 2[-(2A + p^2B)\eta_{\mu\mu'} + 2Bp_\mu p_{\mu'}] , \quad (13)$$

where the scalar quantities A and B are

$$A(p^2, q^2, q'^2) = \frac{1}{(2\pi)^2} \frac{\pi}{24} p^2 \lambda^{3/2}(1, q^2/p^2, q'^2/p^2) ,$$
$$B(p^2, q^2, q'^2) = \frac{1}{(2\pi)^2} \frac{\pi}{12} \lambda^{1/2}(1, q^2/p^2, q'^2/p^2) \left(1 + \frac{q^2}{p^2} + \frac{q'^2}{p^2} - 2\frac{(q^2 - q'^2)^2}{p^4}\right) , \quad (14)$$

with $\lambda(x,y,z,) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ the usual triangular function.

The cross-section can then be written as

$$\left(\frac{d\sigma}{dQ_1^2\,dQ_2^2}\right)^{WW^*} = \frac{1}{2\pi}\,\frac{(g/\sqrt{2})^2 2(Q_1^2/24\pi)}{|D_W(Q_1^2)|^2}\,\frac{1}{2\pi}\,\frac{3\cdot 2(g/\sqrt{2})^2|V_{tb}|^2}{|D_W(Q_2^2)|^2}\,\frac{1}{2s}\int d^2PS(K;p_1,p_2)\widetilde{\mathcal{A}}^2\,, \tag{15}$$

where

$$\widetilde{\mathcal{A}}^2 = \tfrac{1}{4}\,\mathrm{Tr}(\not{k}_2 T_{\mu\nu}\not{k}_1\gamma_0 T^{\dagger}_{\mu'\nu'}\gamma_0)(-\eta^{\nu\nu'} + p_1{}^{\nu}p_1{}^{\nu'}/Q_1{}^2)(-F\eta^{\mu\mu'} + Gp_2^{\mu}p_2^{\mu'}/M_W^2)\,, \tag{16}$$

with

$$F(Q_2^2,m_t^2,m_b^2) = 2A + Q_2^2B\,,\quad G(Q_2^2,m_t^2,m_b^2) = 4A + 2B(M_W^2 - Q_2^2) - \frac{Q_2^2}{M_W^2}(2A - BQ_2^2)\,. \tag{17}$$

and A and B are given by eq. (14) and evaluated at $p^2 = Q_2^2$, $q^2 = m_t^2$ and $q'^2 = m_b^2$.

Using the narrow width approximation for the W-boson decaying into light leptons, the integration over Q_1^2 becomes trivial and we obtain

$$\left(\frac{d\sigma}{dQ_2^2}\right)^{WW^*} = \mathrm{Br}(W^- \to e^-\overline{\nu})\,\frac{1}{\pi}\,\frac{3(g/\sqrt{2})^2|V_{tb}|^2}{|D_W(Q_2^2)|^2}\,\frac{1}{2s}\int d^2PS(K;p_1,p_2)\widetilde{\mathcal{A}}^2\Big|_{Q_1^2=M_W^2}\,. \tag{18}$$

The factor $\widetilde{\mathcal{A}}^2$, in terms of the variables

$$s = (k_1 + k_2)^2 = (p_1 + p_2)^2\,,\quad t = (k_1 - p_1)^2 = (k_2 - p_2)^2\,,\quad u = (k_1 - p_2)^2 = (k_2 - p_1)^2 \tag{19}$$

can be written as

$$\widetilde{\mathcal{A}}^2 = 2e^4\{(a^2 + b^2)[\xi_1\widetilde{A}_1(s,t,u) + \xi_2\widetilde{A}_2(s,t,u)] + 4(a+b)c[\xi_1\widetilde{I}_1(s,t,u) + \xi_2\widetilde{I}_2(s,t,u)\ ]$$
$$+ 8c^2[\xi_1\widetilde{E}_1(s,t,u) + \xi_2\widetilde{E}_2(s,t,u)]\}\,, \tag{20}$$

where we have introduced the functions $\xi_1 = (Q_2^2/M_W^2)G(Q_2^2,m_t^2,m_b^2)$ and $\xi_2 = F(Q_2^2,m_t^2,m_b^2) - \xi_1$ along with the quantities (again in the notation of ref. [7]) $a = -1 + (e_Z g_V/e^2)s/(s - M_Z^2)$, $b = (e_Z g_A/e^2)s/(s - M_Z^2)$ and $c = (g/2\sqrt{2}/e)^2$ in order to make the comparison with the known result of $e^+e^- \to W^+W^-$ easier. For brevity, we only give the functions ($A_{1,2}$, $I_{1,2}$ and $E_{1,2}$), obtained after integration over the phase space:

$$A_1 = \int dt\,\widetilde{A}_1(s,t,u) = \frac{s^3}{Q_1^2Q_2^2}\lambda^{3/2}\left[\frac{1}{24} + \frac{5}{12}\left(\frac{Q_1^2 + Q_2^2}{s}\right) + \frac{1}{24}\left(\frac{Q_1^4 + Q_2^4 + 10Q_1^2Q_2^2}{s^2}\right)\right],$$

$$A_2 = \int dt\,\widetilde{A}_2(s,t,u) = \frac{s^3}{Q_1^2Q_2^2}\lambda^{3/2}\left[\frac{5}{12}\frac{Q_2^2}{s} + \frac{1}{24}\left(\frac{Q_2^4 + 10Q_2^2Q_1^2}{s^2}\right)\right] - s\lambda^{1/2}\left(1 + \frac{Q_1^2}{s} - \frac{Q_2^2}{2s}\right),$$

$$I_1 = \int dt\,\widetilde{I}_1(s,t,u) = \frac{s^3}{Q_1^2Q_2^2}\lambda^{1/2}\left[\frac{1}{24} + \frac{3}{8}\left(\frac{Q_1^2 + Q_2^2}{s}\right) - \frac{1}{24}\left(\frac{9Q_1^4 + 9Q_2^4 + 10Q_1^2Q_2^2}{s^2}\right)\right.$$
$$\left. - \frac{1}{24}\left(\frac{Q_1^6 + Q_2^6 + 11Q_1^2Q_2^4 + 11Q_1^4Q_2^2}{s^3}\right)\right] + \left(Q_1^2 + Q_2^2 + \frac{Q_1^2Q_2^2}{s}\right)\log L\,,$$

$$I_2 = \int dt\,\widetilde{I}_2(s,t,u) = \frac{s^3}{Q_1^2Q_2^2}\lambda^{1/2}\left[\frac{11}{24}\frac{Q_2^2}{s} - \frac{5}{12}\left(\frac{Q_2^2(Q_2^2 + Q_1^2)}{s^2}\right) - \frac{1}{24}Q_2^2\left(\frac{10Q_1^2Q_2^2 + 13Q_1^4 + Q_2^4}{s^3}\right)\right]$$
$$+ \left(Q_1^2 + Q_2^2 + \frac{Q_1^2Q_2^2}{s}\right)\log L\,, \tag{21}$$

$$E_1 = \int \mathrm{d}t\, \widetilde{E}_1(s,t,u) = \frac{s^3}{Q_1^2 Q_2^2} \lambda^{1/2} \left[\frac{1}{24} + \frac{10}{24}\left(\frac{Q_1^2 + Q_2^2}{s} \right) + \left[-\frac{25}{12} + \frac{1}{24}\left(\frac{Q_1^4 + Q_2^4}{Q_1^2 Q_2^2} \right) \right] \frac{Q_1^2 Q_2^2}{s^2} \right]$$
$$- s\left(1 - \frac{Q_1^2}{s} - \frac{Q_2^2}{s} \right) \log L ,$$
$$E_2 = \int \mathrm{d}t\, \widetilde{E}_2(s,t,u) = \frac{s^3 \lambda^{1/2}}{Q_1^2 Q_2^2} \left(\frac{Q_2^2}{2s} - 2\frac{Q_1^2 Q_2^2}{s^2} \right) + (Q_1^2 + Q_2^2 - s) \log L , \qquad (21\ \text{cont'd})$$

where

$$\lambda = \lambda\left(1, \frac{Q_1^2}{s}, \frac{Q_2^2}{s}\right), \quad L = \frac{(1-\lambda^{1/2})^2 - [(Q_1^2 - Q_2^2)/s]^2}{(1+\lambda^{1/2})^2 - [(Q_1^2 - Q_2^2)/s]^2} .$$

The differential cross section as a function of the invariant mass of the off-shell W boson is then written as

$$\left(\frac{\mathrm{d}\sigma}{\mathrm{d}Q_2^2} \right)^{WW^*} = \frac{2\pi\alpha^2}{s^2} \mathrm{Br}(W \to e^- \overline{\nu}) \frac{1}{\pi} \frac{3(g/\sqrt{2})^2 |V_{tb}|^2}{(Q_2^2 - M_W^2)^2 + (M_W \Gamma_W)^2}$$
$$\times [(a^2+b^2)(\xi_1 A_1 + \xi_2 A_2) + 4(a+b)c(\xi_1 I_1 + \xi_2 I_2) + 8c^2(\xi_1 E_1 + \xi_2 E_2)]|_{Q_1^2 = M_W^2} , \qquad (22)$$

in which the dependence on m_t and m_b is through the functions ξ_1 and ξ_2. A check of the above result can easily be done by considering the W^+ decaying into *light quarks*, and using again the narrow width approximation. Then as $Q_2^2 \to M_W^2$, $F \to G$, $\xi_2 \to 0$ and $\xi_1 \to M_W^2/(24\pi)$ one obtains the expected result: $\sigma^{WW^*} = \sigma(e^+e^- \to W^+W^-)\mathrm{Br}(W^- \to e^-\overline{\nu})\mathrm{Br}(W^+ \to q_i \overline{q}_j)$, where $\sigma(e^+e^- \to W^+W^-)$ coincides with that of ref. [7].

The integrated cross-section for the production of a $t\overline{b}$ pair through the WW^* and tt^* channels is given, respectively, by

$$\sigma^{WW^*}(s; m_t^2, m_b^2) = \int_{(m_t+m_b)^2}^{(\sqrt{s}-M_W)^2} \mathrm{d}Q_2^2 \left(\frac{\mathrm{d}\sigma}{\mathrm{d}Q_2^2} \right)^{WW^*} , \quad \sigma^{tt^*}(s; m_t^2, m_b^2) = \int_{m_b^2}^{(\sqrt{s}-m_t)^2} \mathrm{d}Q^2 \left(\frac{\mathrm{d}\sigma}{\mathrm{d}Q^2} \right)^{tt^*} . \qquad (23)$$

Let us now discuss the production of a single top quark through the process

$$e^+(p_a)\,\gamma(p_b) \to \nu(p_1)\,t(p_2)\,\overline{b}(p_3) , \qquad (24)$$

which proceeds through the diagrams of fig. 2. Its integrated cross section convoluted with the probability distribution function of the photon in the initial electron (Wiezsäcker–Williams approximation) [8], can be used to estimate the rate of production of a single top at e^+e^- colliders. We note that this process is also worthwile to study, in view of recent proposals [9,10] regarding future electron–photon colliders with energies up to $\sqrt{s_{\gamma e}} = 1.6$ TeV and $\mathcal{L} = 10^{33}$ cm^{-1} s^{-1}.

In contrast to the case of W–gluon fusion, here there are two additional diagrams figs. 2c–2d. The gauge invariant amplitude reads

$$\mathcal{M} = (g/2\sqrt{2})^2 e\, V_{tb} (T_\lambda^{(1)} + T_\lambda^{(2)} + T_\lambda^{(3)} + T_\lambda^{(4)}) \epsilon^\lambda(p_b) , \qquad (25)$$

where

$$T_\lambda^{(1)} = \frac{e_b}{t_1 - M_W^2} J^\mu_{\text{lepton}} \overline{u}(p_2) \gamma_\mu (1-\gamma_5) \frac{-1}{\not{p}_b - \not{p}_3 - m_3} \gamma_\lambda v(p_3) ,$$
$$T_\lambda^{(2)} = \frac{e_t}{t_1 - M_W^2} J^\mu_{\text{lepton}} \overline{u}(p_2) \gamma_\lambda \frac{-1}{\not{p}_2 - \not{p}_b - m_2} \gamma_\mu (1-\gamma_5) v(p_3) , \qquad (26)$$

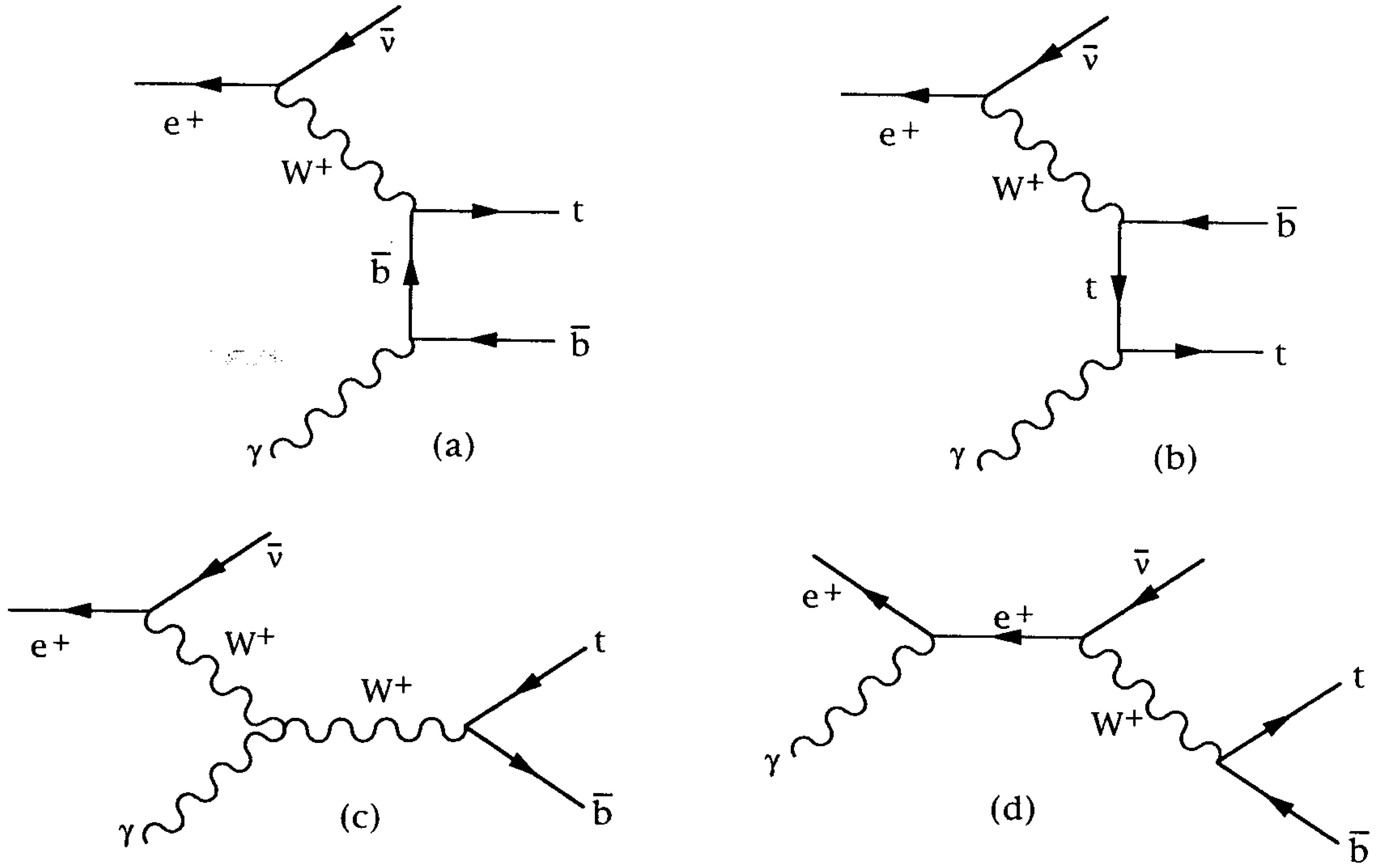

Fig. 2. Lowest order Feynman diagrams for the process $e^+\gamma \to \nu t\bar{b}$.

$$T_\lambda^{(3)} = J^\mu_{\text{lepton}} S_{\mu\nu\lambda}\left(\eta^{\nu\sigma} - \frac{Q_\nu Q_\sigma}{M_W^2}\right) J_\sigma^{\text{quark}} \frac{1}{(t_1 - M_W^2)(s_2 - M_W^2)},$$

$$T_\lambda^{(4)} = \frac{-1}{s(s_2 - M_W^2)} \overline{v}(p_a)\gamma_\lambda(\not{p}_a + \not{p}_b)\gamma^\mu(1-\gamma_5)v(p_1)\left(\eta_{\mu\nu} - \frac{Q_\nu Q_\sigma}{M_W^2}\right) J^\nu_{\text{quark}}, \tag{26 cont'd}$$

with

$$S_{\mu\nu\lambda} = \eta_{\mu\nu}(-Q-P)_\lambda + \eta_{\mu\lambda}(P - p_b)_\nu + \eta_{\nu\lambda}(p_b + Q)_\mu,$$

$$J^\mu_{\text{lepton}} = \overline{v}(p_a)\gamma^\mu(1-\gamma_5)v(p_1), \quad J^\nu_{\text{quark}} = \overline{u}(p_2)\gamma^\nu(1-\gamma_5)v(p_3). \tag{27}$$

After some algebra, it is possible to show that some of the terms in $T_\lambda^{(3)}$ and $T_\lambda^{(4)}$ cancel, so that we can equivalently write

$$T_\lambda^{(3)} = J^\mu_{\text{lepton}} \Gamma_{\mu\lambda\sigma} J^\sigma_{\text{quark}} \frac{1}{(t_1 - M_W^2)(s_2 - M_W^2)},$$

$$T_\lambda^{(4)} = \frac{-1}{s(s_2 - M_W^2)} \overline{v}(p_a)\gamma_\lambda(\not{p}_a + \not{p}_b)\gamma_\mu(1-\gamma_5)v(p_1) J^\mu_{\text{quark}}, \tag{28}$$

where we have defined

$$\Gamma_{\mu\lambda\sigma} = -2[\eta_{\mu\lambda}(p_b)_\sigma + \eta_{\mu\sigma}P_\lambda - \eta_{\lambda\sigma}(p_b)_\mu],$$

$$P = p_a - p_1 \quad \text{and} \quad P^2 = t_1, \qquad Q = p_2 + p_3 \quad \text{and} \quad Q^2 = s_2. \tag{29}$$

Parametrizing the three-particle phase space [11] in the frame $\boldsymbol{p}_2 + \boldsymbol{p}_3 = \boldsymbol{0}$, where each four-vector is written in terms of the two invariants s_2 and t_1 and the polar angles Ω_3 of the vector $\boldsymbol{p}_3$, the cross section may be cast in the form

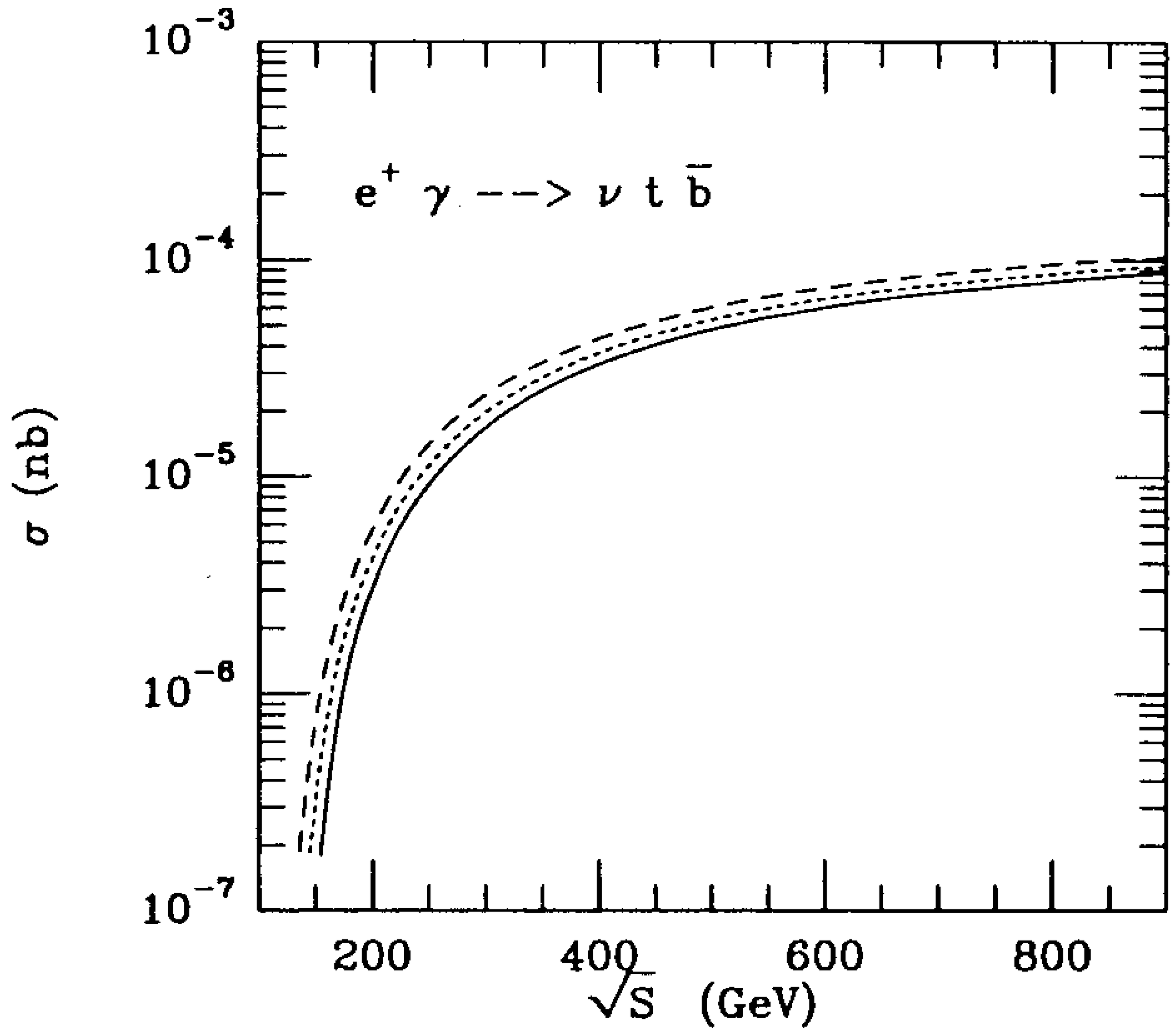

Fig. 3. Integrated cross section for the process $e^+\gamma \to \nu t\bar{b}$ as a function of the center of mass energy. Long dashed curve $m_{top} = 110$ GeV. Small dashed curve $m_{top} = 120$ GeV. Solid curve $m_{top} = 130$ GeV.

$$\widehat{\sigma} = \frac{3\alpha^3}{16\pi x_w^2} \frac{1}{\widehat{s}^2} \int_{(m_t+m_b)^2/\widehat{s}}^{\widehat{s}} ds_2 \int_{-(s-s_2)}^{0} dt_1 \int d\Omega_3 \frac{\lambda^{1/2}(s_2, m_t^2, m_b^2)}{s_2} \times \sum_{i \geqslant j} A_{ij}, \tag{30}$$

with

$$A_{ij} = \frac{1}{64 \cdot 4} \sum_{\text{spins}} \frac{1}{1+\delta_{ij}} [T^i_\lambda (T^j_{\lambda'})^* + T^j_\lambda (T^i_{\lambda'})^*](-\eta^{\lambda\lambda'}). \tag{31}$$

The quantities A_{ij} are cumbersome functions of the particle's momenta scalar products, and are available upon request. The result of the integration is shown in fig. 3. We note that the rate is quite interesting. Above threshold, it produces a substantial number of events (for the $e\gamma$ collider luminosity quoted above).

At e^+e^- colliders, a single top quark production rate arising through the $W^*\gamma^*$ fusion, may be estimated in the equivalent photon approximation as

$$\sigma = \int_{z_{min}}^{1} f_\gamma(z)\,\widehat{\sigma}(zs), \tag{32}$$

with $z_{min} = (m_t + m_b)^2/s$, and where $f_\gamma(x) = (\alpha/\pi)\{[1 + (1-x)^2]/x\}\log(E/m_e)$ is the photon distribution [8] for an electron of energy E and mass m_e.

Numerical estimates of the integrated cross sections for the two different production mechanisms (WW^*, $W^*\gamma^*$) of a single top quark at e^+e^- colliders, and for different center of mass energies of the electron–positron pair, are shown in fig. 4 as a function of the top quark mass, and compared with the production cross-section for a tt^* pair, eg. $e^+e^- \to t\bar{t}^* \to tW^*\bar{b} + e\nu$. In the calculation, we have used $|V_{tb}| = 1.0$, $\mathrm{Br}(W \to e\nu) = 0.11$ and $m_b = 5$ GeV. Our signal cross-sections should be multiplied by a factor of two in order to include the charge conjugate final state (the WW^* and tt^* may also be summed over other light lepton channels).

However, with the presently expected integrated luminosity at LEP II of $\mathcal{L} = 5 \times 10^5$ nb^{-1}/yr, the processes that we have discussed cannot be used to detect the top quark, but, were the luminosity to be ameliorated, their contribution could be not entirely negligible.

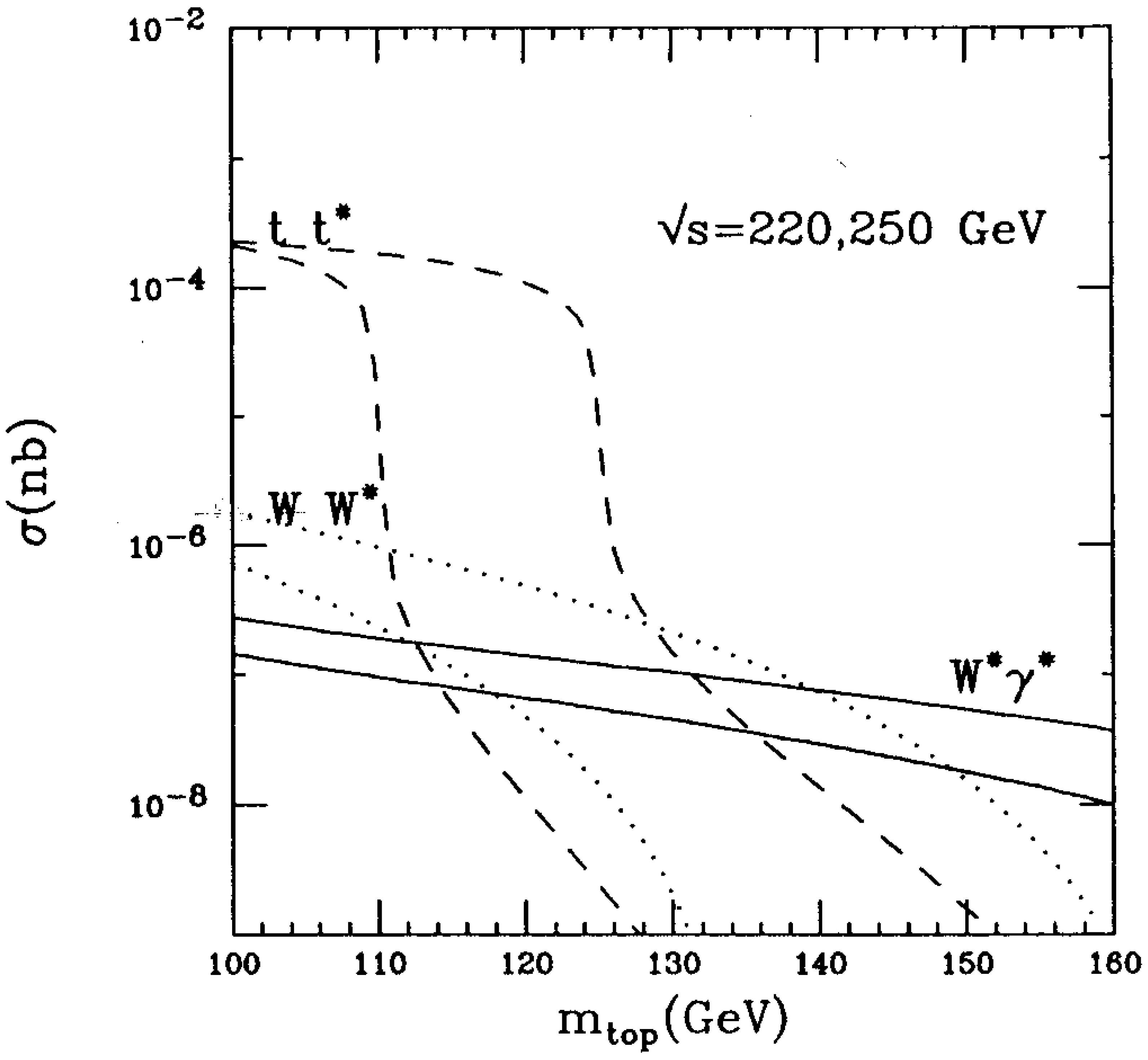

Fig. 4. Integrated cross sections at $\sqrt{s} = 220, 250$ GeV for $e^+e^- \to t\bar{b}e^-\bar{\nu}$. The solid line is the $W^*\gamma^*$ channel, the small dashed curve is the WW^* channel, and the long dashed line is the tt^* channel. Upper curves refer to higher CM energy.

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