Final-state-interaction simulation of T violation in top-quark semileptonic decay

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The standard electroweak final-state interaction induces a false T-odd correlation in the top-quark semileptonic decay. The correlation parameter is calculated in the standard model and found to be considerably larger than those that could be produced by genuine T-violation effects in a large class of theoretical models.

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I. INTRODUCTION

Final-state interactions play an important role in the determination of CP and T violation. A test for CP violation is to compare the partial decay rates of a particle and its antiparticle. In this case final-state interactions are necessary since in their absence the partial decay rates are equal from CPT invariance even if CP is violated. A general formalism for calculating such partial rate differences based on CPT invariance and unitarity has recently been developed [1], and its application to B meson decays (Ref. [1]) and to t-quark decays [2] has revealed some interesting relations between final-state interaction and CP-violation observables in weak decays.

A test [3] for T violation is to observe a "T-odd correlation," such as those of the form $\boldsymbol{\sigma}\cdot(\mathbf{p}_1\times\mathbf{p}_2)$ where $\boldsymbol{\sigma}$ is a spin and $\mathbf{p}_1$ and $\mathbf{p}_2$ are momenta. In contrast with the partial decay difference, a T-odd correlation can be produced by final-state interactions even if T invariance holds. Thus, to use such correlations as a test of T violation the final-state-interaction effect must be negligible or calculable.

This article will be concerned with the t-quark semileptonic decay $t\to bW\to b\nu_l\bar{l}$ in the standard model. Copious production of t quarks at future high-energy colliders such as the Superconducting Super Collider (SSC) and CERN Large Hadron Collider (LHC) have aroused considerable interest in exploring the origin of CP and T violation via t-quark interactions [4]. In particular, a recent study [5] of the possibility of using the T-odd correlation has shown that it has a reasonable sensitivity to some nonstandard sources of T violation. Since such correlations can be produced by standard model physics alone, it is timely to undertake a computation of the final-state-interaction effect due entirely to the standard electroweak interaction, which, up to the one-loop level, respects T and CP invariance in Cabibbo-allowed weak decays such as $t\to bW^+\to b\nu_l\bar{l}$.

II. FINAL-STATE-INTERACTION EFFECT

The computation of final-state-interaction effects on the T-odd correlation has long been of interest. Early examples of the calculation involved nuclear β decay [6], hyperon semileptonic decay [7], and $K_{l3}^{\pm,0}$ decays [8]. The parameter of interest is the coefficient of the T-odd correlation term in the decay spectrum, which in nuclear β decay, for instance, has the following form in the leading approximation:

$$\frac{d\Gamma}{d\Omega_e d\Omega_{\nu_e} dE_e}\sim 1+a\frac{\mathbf{p}_e\cdot\mathbf{p}_{\nu_e}}{E_eE_{\nu_e}} +\boldsymbol{\sigma}\cdot\left[A\frac{\mathbf{p}_e}{E_e}+B\frac{\mathbf{p}_{\nu_e}}{E_{\nu_e}}+D\frac{\mathbf{p}_e\times\mathbf{p}_{\nu_e}}{E_eE_{\nu_e}}\right], \tag{1}$$

where $\boldsymbol{\sigma}$ is the polarization of the parent nucleus and $\mathbf{p}_e(E_e)$ and $\mathbf{p}_{\nu_e}(E_{\nu_e})$ are the electron and neutrino momentum (energy), respectively. In this example, the dominant contribution arises from electromagnetic final-state interaction. The effect depends, among other things, on the recoil of the decaying particle, and thus the size of the T-odd correlation parameter D is of order $D\sim\alpha E_e/M$ $(Z\alpha E_e/M)$ in neutron (nuclear) β decay, where M is a nucleon mass. Since E_e is typically of order 1 MeV, the recoil effect, which is characterized by the ratio E_e/M, is rather tiny. Hence D is highly suppressed in neutron β decay with D typically of the order of $10^{-5}-10^{-6}$. A considerably larger result $(10^{-3}-10^{-4})$ can be obtained in some nuclear β decays due to the enhancement $Z\gg 1$ (Ref. [6]). The typical value of the T-odd correlation is between 10^{-3} and 10^{-4} in a neutral K_{l3}^{0} decay. The result in a charged $K_{l3}^{\pm}$ decay is still smaller $(10^{-5}-10^{-6})$, because there the final-state pion is neutral and the effect can only arise from two-loop graphs.

In terms of weak-current interactions, the t-quark semileptonic decay is analogous in many respects to the nuclear β decay. However, the disparity between m_t and m_b implies that the T-odd correlation in the decay $t\to b\nu_l\bar{l}$ does not have a recoil suppression. Indeed, compared to nuclear β decay, where the recoil effect is of order 10^{-3}, in the t semileptonic decay such effects are given by $E_{\bar{e}}/m_t$, which is of order unity. As a consequence, we expect that the final-state-interaction contribution to the T-odd correlation parameter is roughly

$$D(t\to bW\to b\nu_e\bar{e})\sim\alpha|Q_d|\frac{E_{\bar{e}}}{m_t}\sim\frac{\alpha}{9}\sim 10^{-3}\ , \qquad (2)$$

where $Q_d=-\frac{1}{3}$ is the b-quark charge, and we have taken $E_{\bar{e}}/m_t\sim\frac{1}{3}$.

In what follows we will concentrate on the decay $t\to bW\to b\nu_e\bar{e}$. Insofar as the lepton mass can be ignored, our result holds for the other t-quark semileptonic decays as well.

A large m_t implies that the decay $t\to b\nu_e\bar{e}$ proceeds dominantly through the W resonance. The smallness of the W width ($\Gamma_W/M_W\approx 0.026$) then makes the calculation of the leading final-state-interaction effect very simple. Neglecting the b-quark and lepton masses, the leading contributions are generated by graphs displayed in Fig. 1 with

$$M(t\to b\nu_e\bar{e}) = \left[\frac{ig}{\sqrt{2}}\right]^2\frac{[\bar{u}_{\nu_e}(p_{\nu_e})\gamma_\lambda L v_{\bar{e}}(p_{\bar{e}})][\bar{u}_b(p')\Gamma^\lambda u_t(p)]}{k^2-M_W^2+i\Gamma_W M_W}\ , \qquad (3)$$

where $k=p-p'$ is the momentum transfer carried by the W, L and R are the helicity projection operators, and the effective vertex Γ^λ, which includes one-loop interaction corrections from Figs. 1(b) and 1(c), can be parametrized as

$$\Gamma^\lambda=F_1(k^2)\gamma^\lambda L-iF_2(k^2)m_t\sigma^{\lambda\mu}k_\mu R\ , \qquad (4)$$

where $\sigma^{\lambda\mu}=(i/2)[\gamma^\lambda,\gamma^\mu]$. Terms of the form $\gamma^\lambda R$ and $\sigma^{\lambda\mu}k_\mu L$ vanish in the limit $m_b=0$. Also, the k^λ term drops out for $m_e=m_{\nu_e}=0$. While the form factor $F_1=1+O(\alpha/\pi)$ introduces a correction to the weak interaction charge g, F_2 gives an anomalous moment to the $\bar{b}tW$ vertex.

In analogous to nuclear β decay one may define a T-odd correlation parameter D:

$$\frac{d\Gamma}{d\Omega}=\frac{g^4}{(2\pi)^5}\frac{m_t E_{\nu_e}E_{\bar{e}}}{|k^2-M_W^2+i\Gamma_W M_W|^2}\times\left[\left[1-\frac{k^2}{2m_t E_{\nu_e}}\right]+D\left[1-\frac{2E_{\bar{e}}}{m_t}\right]\boldsymbol{\sigma}_t\cdot\frac{\mathbf{p}_{\bar{e}}\times\mathbf{p}_{\nu_e}}{E_{\bar{e}}E_{\nu_e}}\right]+\cdots \qquad (5)$$

with

$$D=m_t^2\,\mathrm{Im}F_2(M_W^2) \qquad (6a)$$

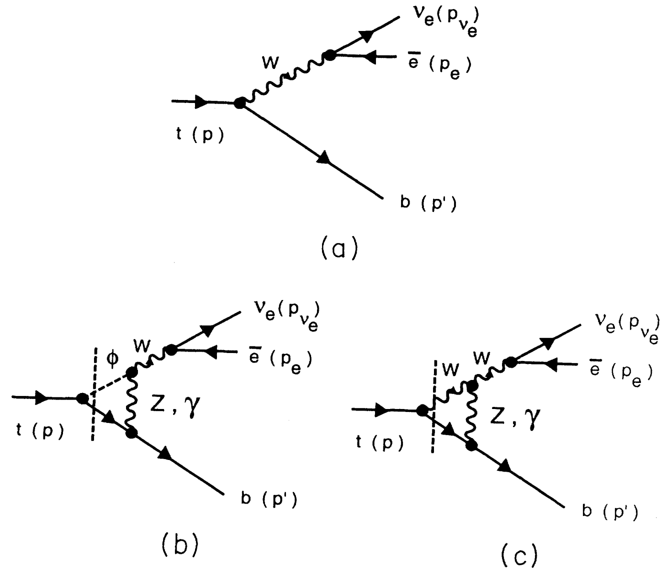

FIG. 1. Feynman graphs generating the dominant contributions to the T-odd correlation. The calculation is carried out in the Feynman–'t Hooft gauge. ϕ is the Higgs-Goldstone boson.

evaluated at $k^2=M_W^2$. The ellipses in Eq. (5) refer to the other terms of no interest to us and $d\Omega=(d^3\mathbf{p}_{\bar{e}}/2E_{\bar{e}})(d^3\mathbf{p}_{\nu_e}/2E_{\nu_e})(d^3\mathbf{p}'/2p_0')$. In reaching (6a) we have taken $F_1=1$.

The final-state interaction in nuclear β decay takes place between the daughter nucleus and the electron. By contrast, the dominant effect in the decay $t\to bW^+\to b\bar{e}\nu_e$ arises from $bW\to bW$ rescattering. By employing the unitarity formula given by Wolfenstein (Ref. [1]) one can show that the relevant interactions are those which scatter a bW^+ state to other bW^+ states with different spin configurations. As a result, the T-odd correlation parameter is directly proportional to the absorptive part of the form factor F_2 which connects hadron states with different helicities. We find (the detail of the calculation is summarized in the Appendix)

$$\mathrm{Im}F_2(M_W^2)=-\frac{\alpha Q_d}{2m_t^2}\left[1-\frac{1}{2}\frac{M_W^2}{m_t^2}\right]+\frac{\alpha(1+2Q_d s^2)}{8(m_t^2-M_W^2)}\left[\left[-\frac{1}{c^2}-\frac{1}{s^2}\right]I_1+\frac{2}{s^2}I_2\right], \qquad (6b)$$

where $s^2=\sin^2\theta_W$, $c^2=\cos^2\theta_W$, and

$$I_1=2+\left[1+\frac{2m_t^2M_Z^2}{(m_t^2-M_W^2)^2}\right]\ln\frac{M_Z^2m_t^2}{M_Z^2m_t^2+(m_t^2-M_W^2)^2}\ ,$$

$$I_2=\left[1-\frac{M_W^2}{m_t^2}\right]\left[1-\frac{1}{2}\frac{M_W^2}{m_t^2}+2\frac{M_Z^2}{m_t^2-M_W^2}+3\frac{M_W^2M_Z^2}{(m_t^2-M_W^2)^2}\right]+\frac{M_Z^2}{m_t^2-M_W^2}\left[2+2\frac{M_Z^2}{m_t^2-M_W^2}+3\frac{M_W^2M_Z^2}{(m_t^2-M_W^2)^2}\right]\ln\frac{M_Z^2m_t^2}{M_Z^2m_t^2+(m_t^2-M_W^2)^2}\ . \qquad (6c)$$

In Eq. (6b) the first term comes from the photon graphs and the second from the Z. For a very heavy top the result is dominated by the Z exchange diagram and has a logarithmic dependence on m_t. Asymptotically it approaches

$$\lim_{m_t \to \infty} D \sim \frac{\alpha}{6} \left\{ 1 - \frac{3}{4} \left[1 - \frac{2s^2}{3} \right] \times \left[\frac{2}{c^2} + \left[\frac{1}{c^2} + \frac{1}{s^2} \right] \ln \frac{M_Z^2}{m_t^2} \right] \right\} . \tag{7}$$

The numerical results for D from Eqs. (6a) to (6c) are summarized [9] in Table I for m_t between 100 and 200 GeV. One sees that D is between 1×10^{-3} and 6×10^{-3}, as we expected from the simple dimensional argument Eq. (2). The result shows a slow increase with larger values of m_t in this region.

The T-odd correlation may be reparametrized in terms of an asymmetry parameter A, which is related to the difference of the decay $W^+ \to \bar{e}\nu_e$ occurring in the opposite sides of the $\sigma_t \times \mathbf{p}$ plane (Ref. [5]):

$$A = -\frac{3(m_t^2 - M_W^2)}{4(m_t^2 + 2M_W^2)} \frac{m_t M_W \mathrm{Im}(F_1 F_2^*)}{|F_1|^2} \approx -\frac{3(m_t^2 - M_W^2)}{4(m_t^2 + 2M_W^2)} m_t M_W \mathrm{Im} F_2^* , \tag{8}$$

where $\mathrm{Im}F_2^*$ is given by Eqs. (6b) and (6c) with an additional overall minus sign. The results for A are summarized in the last column of Table I. They vary from 1×10^{-4} to 1×10^{-3} for $m_t = 100-200$ GeV. In comparison with the maximal-allowed T violation effect in the models considered in Ref. [5] in which $A < 5\times10^{-5} \sim 5\times10^{-4}$, the standard model final-state interaction produces a much larger false effect.

It is difficult to calculate the T-odd parameter to an accuracy of $\sim 30\%$. The major theoretical uncertainties of the present calculation come from neglecting QCD corrections, which introduce a sizable interference between the absorptive part of F_1 from electroweak interactions and the real part of F_2 from QCD. An order of $\sim(1\sim10)\%$ correction due to this effect alone is possible. A still more complicated contribution arises from the interference between $\mathrm{Im}F_2$ calculated above and the real part of F_1 due to QCD. Other uncertainties arise from neglecting (1) the WZ threshold effect (relevant if $m_t > M_W + M_Z + m_b$) and (2) all the box diagrams. The contribution of the latter also depends on the angle between $\mathbf{p}_{\bar{e}}$ and $\mathbf{p}'$ in a rather complicated way. All of these contributions are suppressed by the ratio Γ_W / M_W, however. The calculation of these next-leading terms would be crucial should future experiments approach the precision of $D \sim 10^{-3}$.

T-odd correlations of the form $\sigma_{\bar{e}} \cdot (\mathbf{p}_{\nu_e} \times \mathbf{p}_{\bar{e}})$, $\sigma_b \cdot (\mathbf{p}_{\nu_e} \times \mathbf{p}_{\bar{e}})$ and P- and CP-odd correlations of the form $\sigma_t \cdot (\sigma_b \times \mathbf{p}')$ are much more difficult to measure experimentally, and thus will not be considered in this paper.

TABLE I. The result for the T-odd correlation in the t-quark semileptonic decay. The parameters D and A are defined in Eqs. (5) and (9), respectively.

m_t (GeV)	D	A
100	1.1×10^{-3}	1.0×10^{-4}
110	1.4×10^{-3}	1.8×10^{-4}
120	1.9×10^{-3}	2.8×10^{-4}
130	2.4×10^{-3}	3.9×10^{-4}
140	2.9×10^{-3}	5.1×10^{-4}
150	3.5×10^{-3}	6.4×10^{-4}
160	4.0×10^{-3}	7.6×10^{-4}
170	4.6×10^{-3}	8.7×10^{-4}
180	5.1×10^{-3}	9.8×10^{-4}
190	5.6×10^{-3}	1.1×10^{-3}
200	6.1×10^{-3}	1.2×10^{-3}

III. CONCLUSION

We have calculated the T-odd correlation $\sigma_t \cdot (\mathbf{p}_{\bar{e}} \times \mathbf{p}_{\nu_e})$ induced by the standard electroweak final-state interactions in the decay $t \to bW^+ \to b\bar{l}\nu_l$, and found that the result has a logarithmic dependence on the t-quark mass and is dominated by the $bW \to bW$ rescattering due to a Z exchange in the heavy top limit. For m_t in the range 100–200 GeV the correlation parameter D defined in Eq. (5) is between 1×10^{-3} and 6×10^{-3}, and the asymmetry parameter A given by Eq. (8) is between 1×10^{-4} and 1×10^{-3}. It is shown that the standard model physics can simulate a false T-odd signal, with its magnitude exceeding genuine T-violation effects of a size that could possibly be produced in a large class of theoretical models. To get rid of this pure final-state interaction effect one may consider comparing the asymmetry parameter for both $t \to bW^+$ and $\bar{t} \to \bar{b}W^-$, as in the study of CP-violating parameters $\alpha + \bar{\alpha}$ and $\beta + \bar{\beta}$ in the Λ decays [10].

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APPENDIX

We give some details of the calculation in this appendix. The technique is standard [11] except that we use the Minkowskian metric $g^{\lambda\beta} = \mathrm{diag}(1, -1, -1, -1)$. The one- and two-point functions are defined as

$$A(m)=-i\mu_0^{(n-4)}\int\frac{d^nK}{(2\pi)^n}\frac{1}{K^2-m^2+i\epsilon}\,,$$
$$B(m_1,m_2;k)=-i\mu_0^{(n-4)}\int\frac{d^nK}{(2\pi)^n}\frac{1}{[K^2-m_1^2+i\epsilon][(K+k)^2-m_2^2+i\epsilon]}\,,\tag{A1}$$

where $\epsilon\to 0_+$, and we use dimensional regularization to isolate the ultraviolet divergences. The only relevant three-point function is

$$C_0=-i\int\frac{d^4K}{(2\pi)^4}\frac{1}{[K^2-M_Z^2+i\epsilon][(K-k)^2-M_W^2+i\epsilon][(K+p')^2-m_b^2+i\epsilon]}\,.\tag{A2}$$

We find

$$\mathrm{Im}A(m)=0\,,$$
$$\mathrm{Im}B(m_1,m_2;k)=\frac{1}{16\pi k^2}\sqrt{\lambda(k^2,m_1^2,m_2^2)}\,\theta[k^2-(m_1+m_2)^2]\,,\tag{A3}$$
$$\mathrm{Im}C_0=\frac{1}{16\pi\sqrt{\lambda(m_t^2,M_W^2,m_b^2)}}\ln\frac{M_Z^2m_t^2}{M_Z^2m_t^2+\lambda(m_t^2,M_W^2,m_b^2)}\,\theta[m_t^2-(M_W+m_b)^2]\,,$$

where $\lambda(x,y,z)=x^2+y^2+z^2-2xy-2xz-2yz$. In evaluating $\mathrm{Im}C_0$ we have put all the external lines on their mass shell.

Neglecting the b-quark and lepton masses, the final-state-interaction effect due to a photon exchange is

$$\Gamma^\lambda(\gamma)=-e^2Q_d4m_tp'^\lambda R(a_1+b_1)\,,\tag{A4}$$

where a_1 and b_1 are the coefficients defined in the integrals

$$-i\int\frac{d^4K}{(2\pi)^4}\frac{K^\lambda}{K^2[(K-k)^2-M_W^2][(K+p')^2-m_b^2]}=a_1k^\lambda+a_2p'^\lambda\,,\tag{A5}$$

$$-i\mu_0^{(n-4)}\int\frac{d^nK}{(2\pi)^n}\frac{K^\lambda K^\beta}{K^2[(K-k)^2-M_W^2][(K+p')^2-m_b^2]}=b_1(k^\lambda p'^\beta+k^\beta p'^\lambda)+b_2g^{\lambda\beta}+b_3k^\lambda k^\beta+b_4p'^\lambda p'^\beta\,.\tag{A6}$$

We find

$$a_1=\frac{B(M_W,0;k)-B(M_W,0;p)}{m_t^2-M_W^2}\,,$$
$$b_1=-\frac{1}{2}\frac{A(M_W)-M_W^2B(0,M_W;p)}{m_t^2(m_t^2-M_W^2)}\,.\tag{A7}$$

It then follows from Eqs. (A3), (A4), and (A7) that the absorptive part of $\Gamma^\lambda(\gamma)$ is

$$\Gamma^\lambda_{\mathrm{abs}}(\gamma)=\frac{\alpha Q_d}{m_t}\left[1-\frac{M_W^2}{2m_t^2}\right]p'^\lambda R\,.\tag{A8}$$

It can be written in a more conventional form by applying the Gordon identity

$$[\bar u_b(p')p'^\lambda Ru_t(p)]=\frac{i}{2}[\bar u_b(p')\sigma^{\lambda\mu}k_\mu Ru_t(p)]+\cdots\,.\tag{A9}$$

The result due to Z exchange is

$$\Gamma^\lambda(Z)=-e^2(1+2Q_ds^2)m_tp'^\lambda R\times\left[\left[-\frac{1}{c^2}-\frac{1}{s^2}\right](-a_1'+a_2'+C_0)-\frac{2}{s^2}(a_1'+b_1')\right]\,,\tag{A10}$$

where the coefficients a_1', a_2', and b_1' are defined analogously:

$$-i\int\frac{d^4K}{(2\pi)^4}\frac{K^\lambda}{[K^2-M_Z^2][(K-k)^2-M_W^2][(K+p')^2-m_b^2]}=a_1'k^\lambda+a_2'p'^\lambda\,,\tag{A11}$$
$$-i\mu_0^{(n-4)}\int\frac{d^nK}{(2\pi)^n}\times\frac{K^\lambda K^\beta}{[K^2-M_Z^2][(K-k)^2-M_W^2][(K+p')^2-m_b^2]}=b_1'(k^\lambda p'^\beta+k^\beta p'^\lambda)+b_2'g^{\lambda\beta}+b_3'k^\lambda k^\beta+b_4'p'^\lambda p'^\beta\,.$$

We find

$$a_1'=\frac{1}{m_t^2-M_W^2}[B(M_Z,M_W;k)-B(M_W,0;p)-M_Z^2C_0]\,,\tag{A12}$$

$$a_2'=-\frac{1}{m_t^2-M_W^2}[B(M_Z,0;p')-B(M_W,0;p)-M_Z^2C_0]-\frac{2M_W^2}{(m_t^2-M_W^2)^2}[B(M_Z,M_W;k)-B(M_W,0;p)-M_Z^2C_0]\,,\tag{A13}$$

$$b_1' = -\frac{1}{(m_t^2-M_W^2)^2}\left\{\frac{1}{2}\left[1-\frac{M_W^2}{m_t^2}\right][A(M_W)-M_W^2B(M_W,0;p)]+\tfrac{1}{2}[A(M_Z)-M_Z^2B(M_W,M_Z;k)]\right.$$
$$+2M_Z^2[B(M_W,0;p)-B(M_Z,0;p')]-3\frac{M_W^2M_Z^2}{m_t^2-M_W^2}[B(M_W,M_Z;k)-B(M_W,0;p)]$$
$$\left.+\left[2+3\frac{M_W^2}{m_t^2-M_W^2}+\frac{m_t^2-M_W^2}{M_Z^2}\right]M_Z^4C_0\right\}. \quad \text{(A14)}$$

One can check that in the limit $M_Z=0$, a_1 and a_1' become identical and so do b_1 and b_1'. The logarithmic dependence on m_t in the limit $m_t\to\infty$ arises because $\Gamma^\lambda(Z)$ has a term which is directly proportional to C_0 [see (A10)].

It then follows that the absorptive part of $\Gamma^\lambda(Z)$ is

$$\Gamma^\lambda_{\rm abs}(Z) = -\frac{\alpha(1+2Q_ds^2)}{4(m_t^2-M_W^2)}m_tp'^\lambda R\left[\left[-\frac{1}{c^2}-\frac{1}{s^2}\right]I_1+\frac{2}{s^2}I_2\right], \quad \text{(A15)}$$

where $I_{1,2}$ are given by Eq. (6c). Adding (A8) and (A15) we obtain the results given by Eqs. (6a)–(6c) of the text.

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