Consequences of a possible $\gamma\gamma$ resonance at KEK TRISTAN

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If high mass $\gamma\gamma$ events observed at CERN LEP are due to the production of a $\gamma\gamma$ resonance via its leptonic coupling, its consequences can be observed at KEK TRISTAN. We find that a predicted Z decay branching rate is too small to account for the observed events if the resonance spin is zero, due to a strong cancellation in the decay amplitudes. Such a cancellation is absent if the resonance has a spin two. We study the consequences of a tensor production in the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$ at energies reached at TRISTAN. Complete helicity amplitudes with tensor boson exchange contributions are given, and the signal can clearly be identified from various distributions. TRISTAN experiments are also sensitive to the virtual tensor boson exchange effects, which reduce to the contact interaction terms in the high mass limit.

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I. INTRODUCTION

Following the observation by the L3 Collaboration of a cluster of events with a 60 GeV $\gamma\gamma$ and a charged lepton pair [1], similar events have been looked for by the other experimental group at the CERN e^+e^- collider LEP [2]. Although the possibility exists that these events can be explained away as a statistical fluctuation of normal radiative Z decays, it is worthwhile to study the consequences of a narrow $\gamma\gamma$ resonance of mass about 60 GeV [3,4].

When viewed as Z decays into a $\gamma\gamma$ resonance, the reported events [1,2] seem to indicate that it is produced only in association with a charged lepton pair, e^+e^- or $\mu^+\mu^-$. In particular, no apparent clustering of events with a $\gamma\gamma$ invariant mass at around 60 GeV has been found in the other channels $\gamma\gamma\nu\bar{\nu}$ and $\gamma\gamma q\bar{q}$. This seems to exclude the possibility that a $\gamma\gamma$ resonance is produced in association with a virtual Z boson, as expected for a Higgs-like boson [3].

One possible explanation is that a $\gamma\gamma$ resonance is produced in Z decays via its direct coupling to charged leptons. With a significant coupling to electrons, the resonance can be produced at KEK TRISTAN and affects the cross sections for the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, $\gamma\gamma$ at e^+e^- center-of-mass energy $\sqrt{s} \simeq 60$ GeV. Searches for such events have been recently carried out at TRISTAN and the absence of any deviation from the standard model (SM) expectation has been reported [5]. The resulting preliminary bounds on the resonance parameters have been reported [5] by assuming that it has spin zero [6].

In this paper, we report that the production of a massive spinless boson, whether it is a scalar or a pseudoscalar or their doublet, via its coupling to charged leptons is strongly suppressed in Z decays due to a cancellation of amplitudes for the almost axial-vector-like Z coupling to charged leptons. The suppression is so strong that a 60 GeV spinless boson should necessarily have as large a width as its mass in order to gain a Z branching fraction of the order of 10^{-6}.

We find on the other hand that such a cancellation does not take place for the production of a tensor (spin-2) boson in Z decays and that a narrow 60 GeV tensor boson can be produced via its leptonic coupling in Z decays with a significant branching fraction. We therefore study the consequences of a massive tensor boson that couples to charged leptons and two photons at e^+e^- collider energies. Helicity amplitudes are given for the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$ such that the differential cross sections for arbitrarily polarized e^+e^- beams are obtained easily. The angular distributions are found to be distinctive for a spin-2 boson exchange. In contrast with the spin-0 boson exchange cases [6], we find significant interference effects between the spin-2 boson exchange amplitudes and the SM ones, due to the chirality-conserving nature of the tensor boson-lepton coupling. The interference effects allow low-energy experiments to have a good sensitivity for the exchange of a very heavy spin-2 boson. We also compare the effects of a tensor boson exchange in the heavy mass limit and those of the contact four-fermion [7] or two-fermion–two-gamma [8] interactions.

The paper is organized as follows. In Sec. II, we study Z decays into a massive boson and a charged lepton pair for the spin-0 and spin-2 cases. The boson-leptonic widths are expressed in terms of the Z decay branching fraction. In Sec. III, the consequences of a massive spin-2 boson exchange in the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$ are studied in detail. The complete helicity amplitudes and the most general differential cross sections for these processes are presented in this section. In Sec. IV, we compare the effects of a tensor boson exchange in the heavy mass limit with those of contact four-fermion and two-fermion–two-gamma interactions. Section V summarizes our findings.

II. Z DECAYS INTO A BOSON AND A CHARGED LEPTON PAIR

In this section, we study the decay

$$Z \longrightarrow X\ell\bar{\ell}$$

via boson X couplings to charged leptons $\ell\bar{\ell}$ ($\ell = e$ or μ). The boson X has a mass around 60 GeV and decays subsequently into two photons, and hence the X spin cannot be one due to Yang's theorem. We study the two simplest cases spin 0 and 2.

A. Spin-0 boson

A doublet of spinless bosons, a scalar ϕ_S, and a pseudoscalar ϕ_P can have significant couplings to light leptons without violating the leptonic chiral invariance [9]. Although their couplings to two photons should necessarily violate the chiral invariance [11], stringent limits on their leptonic couplings from the electron and muon anomalous magnetic moment measurements can be avoided in the one-loop order if the doublet masses are almost degenerate [10].

We adopt the following effective Lagrangian [6] for the spinless boson doublet and the charged leptonic fields ψ_ℓ:

$$\mathcal{L} = \overline{\psi_\ell}\,(f_S\phi_S + if_P\gamma_5\,\phi_P)\,\psi_\ell + \text{H.c.}\,, \tag{2.1}$$

for $\ell = e$ and μ. The leptonic chiral invariance is preserved in the above interaction if the doublet masses are degenerate ($m_S = m_P$) and the couplings are the same ($f_S = f_P$). The Z boson can decay into a boson and a charged lepton pair via the two Feynman diagrams of Fig. 1. It is instructive to study the helicity amplitude of the decay in the massless lepton limit. If the final lepton

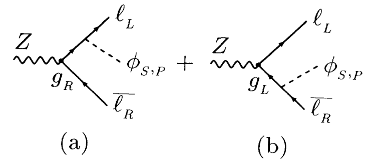


FIG. 1. Feynman diagrams for the decay $Z \to \phi_{S,P}\ell_L\overline{\ell_R}$.

is left handed, the Z boson couples to the right-handed lepton in Fig. 1(a) because of the helicity flip nature of the spinless boson coupling (2.1), whereas in Fig. 1(b) the Z boson couples to the left-handed lepton. If we denote the Z boson lepton coupling as

$$\mathcal{L} = \sum_l \overline{\psi_\ell}\,\gamma^\mu\,(g_L P_L + g_R P_R)\,\psi_\ell Z_\mu\,, \tag{2.2}$$

with the chiral projectors $P_L = (1-\gamma_5)/2$ and $P_R = (1+\gamma_5)/2$, the SM couplings are

$$\begin{aligned} g_L &= g_Z(-\tfrac{1}{2} + \sin^2\theta_W)\,, \\ g_R &= g_Z \sin^2\theta_W\,, \end{aligned} \tag{2.3}$$

with $g_Z = g/\cos\theta_W = e/\sin\theta_W\cos\theta_W$. For $\sin^2\theta_W = 0.23$, the two couplings have almost the same magnitude and opposite signs. We find that, for a massive spinless boson, the two amplitudes tend to cancel because of this near cancellation of the Z boson leptonic couplings in Eq. (2.3).

More explicitly, the helicity amplitudes for the process

$$Z(q,\lambda) \longrightarrow \phi_{S,P}(k) + \ell\,(p,\sigma) + \bar{\ell}\,(\bar{p},\bar{\sigma}) \tag{2.4}$$

are nonvanishing for massless leptons only when the two lepton helicities agree ($\sigma = \bar{\sigma}$), and they are given as

$$M(\lambda,\sigma=L) = g_R\bar{u}(p,L)\,(f_S + if_P\gamma_5)\,\frac{1}{\not{p}+\not{k}}\,\not{\epsilon}(q,\lambda)\,v(\bar{p},L) + g_L\bar{u}(p,L)\,\not{\epsilon}(q,\lambda)\frac{-1}{\not{\bar{p}}+\not{k}}(f_S + if_P\gamma_5)\,v(\bar{p},L) \tag{2.5}$$

for left-handed leptons ($\sigma = L$). The amplitude for right-handed lepton ($\sigma = R$) production is obtained from (2.5) simply by replacing the index L with R. The cancellation between the two terms of Eq. (2.5) can be shown as follows. We express the squared amplitudes summed over helicities as

$$\sum_{\lambda,\sigma}|M(\lambda,\sigma)|^2 = (f_S^2 + f_P^2)\left\{(g_L+g_R)^2\,\Sigma_V(s_1,s_2) + (g_L-g_R)^2\,\Sigma_A(s_1,s_2)\right\}\,, \tag{2.6}$$

and find

$$\Sigma_V = (s_1s_2 - m^2s)\left(\frac{1}{s_1^2}+\frac{1}{s_2^2}\right) + \frac{2(s_2-m^2)(s_1-m^2)}{s_1s_2}\,, \tag{2.7}$$

$$\Sigma_A = (s_1-m^2)(s_2-m^2)\left(\frac{1}{s_1}-\frac{1}{s_2}\right)^2 + \left(\frac{2}{s}-\frac{m^2}{s_1^2}-\frac{m^2}{s_2^2}\right)(s-s_1-s_2+m^2)\,, \tag{2.8}$$

for the degenerate mass case ($m_S = m_P = m$) where $s_1 = (k+p)^2$ and $s_2 = (k+\bar{p})^2$ are Dalitz variables and $s = m_Z^2$. We show in Figs. 2(a) and 2(b) the Dalitz distributions Σ_V and Σ_A, respectively. Σ_V is the term which should be dominant if the Z couplings to charged leptons were vectorlike, whereas Σ_A dominates in the SM where the Z couplings to charged leptons are almost axial-vector-like. The cancellation between the two amplitudes in

Σ_A is clearly seen in Fig. 2(b) along the symmetric line $s_1 = s_2$.

The Z boson decay distribution is expressed as

$$d\Gamma = \frac{1}{2m_z}\frac{1}{3}\sum_{\lambda,\sigma}|M(\lambda,\sigma)|^2\, d\Phi\,, \tag{2.9}$$

with the phase space factor

$$d\Phi = \frac{1}{128\pi^3 m_z^2}\, ds_1 ds_2 \tag{2.10}$$

parametrized in terms of the Dalitz variables. Upon integration over the phase space, we find, for $m_s = m_p = m$,

$$\Gamma(Z\to\phi\,\ell\,\bar{\ell}) = \frac{(f_s^2+f_p^2)\,m_z}{768\,\pi^3}\left[(g_L+g_R)^2 F_V(m^2/m_z^2) + (g_L-g_R)^2 F_A(m^2/m_z^2)\right], \tag{2.11}$$

where

$$F_V(x) = -2+8x-6x^2-(1-2x-3x^2)\ln x - x^2\left\{4\,\mathrm{Li}_2\left(\frac{1}{1+x}\right)-\frac{\pi^2}{3}+2\ln^2(1+x)-\ln^2 x\right\}, \tag{2.12}$$

$$F_A(x) = -\frac{11}{3}-5x+9x^2-\frac{1}{3}x^3-(1+8x+3x^2)\ln x + x^2\left\{4\mathrm{Li}_2\left(\frac{1}{1+x}\right)-\frac{\pi^2}{3}+2\ln^2(1+x)-\ln^2 x\right\}. \tag{2.13}$$

Here Li_2 is the dilogarithm $\mathrm{Li}_2(z) = -\int_0^z (dt/t)\ln(1-t)$. At $x = (59\text{ GeV}/91.187\text{ GeV})^2 = 0.419$, we find $F_V(x) = 0.0186$ and $F_A(x) = 0.000\,619$. The boson-leptonic width is

$$\Gamma(\phi\to\ell\,\bar{\ell}) = \frac{(f_s^2+f_p^2)\,m}{8\,\pi} \tag{2.14}$$

for a degenerate doublet. We therefore find from Eqs. (2.11)–(2.14) that

$$\Gamma(\phi\to\ell\,\bar{\ell}) = (1.56\times10^7\text{ GeV})B(Z\to\phi\,\ell\,\bar{\ell}) \tag{2.15}$$

for $m = 59$ GeV, $\sin^2\theta_W = 0.23$, and $g_z^2 = 4\pi\alpha(m_z^2)/\sin^2\theta_W\cos^2\theta_W$ with $\alpha(m_z^2) = 1/128$. If we require that the branching fraction $\sum_{\ell=e,\mu}B(Z\to\phi\,\ell\,\bar{\ell})$ to be larger than 10^{-6}, then the leptonic width $\sum_{\ell=e,\mu}\Gamma(\phi\to\ell\,\bar{\ell})$ should be larger than 15 GeV. This is clearly incompatible with the narrow width assumption of the $\gamma\gamma$ resonance. It is clear that no improvement of the situation is possible by introducing a splitting in the scalar and the pseudoscalar masses.

B. Spin-2 boson

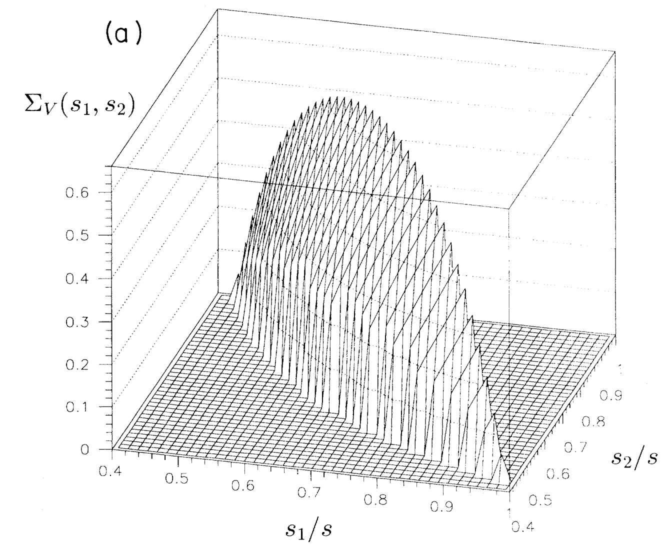
We therefore examine production of a spin-2 (tensor) boson in Z decays:

$$Z(q,\lambda) \longrightarrow \phi_T(k,\tau) + \ell(p,\sigma) + \bar{\ell}(\bar{p},\bar{\sigma})\,, \tag{2.16}$$

where the indices λ, τ, σ, and $\bar{\sigma}$ denote helicities of Z, the tensor ϕ_T, the lepton, and the antilepton, respectively. The chirality-conserving couplings of a tensor boson to a lepton can be parametrized as

$$\mathcal{L} = \frac{i}{2}\phi_T^{\mu\nu}\sum_{\alpha=L,R} f_\alpha^\ell\left[\overline{\psi_{\ell\alpha}}\gamma_\mu(D_\nu\psi_{\ell\alpha}) - (\overline{D_\nu\psi_{\ell\alpha}})\gamma_\mu\psi_{\ell\alpha}\right], \tag{2.17}$$

where D_ν denotes the gauge-covariant derivative of the



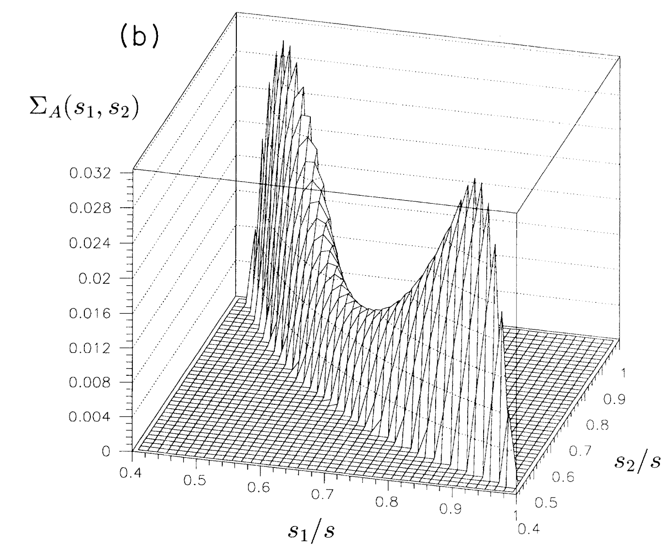


FIG. 2. Dalitz distributions (a) Σ_V and (b) Σ_A of Eqs. (2.7) and (2.8), respectively, for $m = 59$ GeV.

SM. The two terms in the square brackets are shown explicitly so that the Lagrangian is Hermitian even for an off-shell tensor boson, whose effects will be studied in the following section. It should be noted that the SM gauge invariance implies that the tensor boson ϕ_T is a singlet when it couples to right-handed charged leptons ($f_R^\ell \neq 0, f_L^\ell = 0$), whereas it is a singlet or a member of a triplet when it couples to left-handed charged leptons ($f_L^\ell \neq 0, f_R^\ell = 0$). The two couplings f_L^ℓ and f_R^ℓ can in general coexist. However, nonobservation of $\gamma\gamma\nu\bar\nu$-type events disfavors scenarios where ϕ_T couples to left-handed leptons. We will hence consider the case with $f_L^\ell = 0$, although we give expressions for general couplings below.

The Z boson decay (2.16) proceeds via the Feynman diagrams of Fig. 3. In contrast with the case of the spinless boson, the helicity amplitudes are nonvanishing when $\sigma = -\bar\sigma$ in the massless lepton limit, and all three diagrams of Fig. 3 give a contribution proportional to the same Z boson leptonic coupling for each helicity amplitude. The amplitudes for the right-handed lepton ($\sigma = R$) can be written as

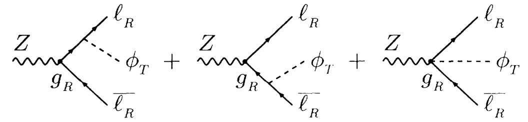


FIG. 3. Feynman diagrams for the decay $Z \to \phi_T \ell_R \overline{\ell_R}$.

$$M(\lambda,\tau,\sigma=R) = g_R f_R^\ell\, \epsilon_T^{\mu\nu}(k,\tau)^*\, \epsilon^\rho(q,\lambda)\bar{u}(p,R) \left\{ \frac{p_\nu}{s_1}\gamma_\mu(\not{p}+\not{k})\gamma_\rho + \frac{\bar{p}_\nu}{s_2}\gamma_\rho(\not{\bar{p}}+\not{k})\gamma_\mu - \gamma_\mu g_{\nu\rho} \right\} P_R v(\bar{p},L)\,, \tag{2.18}$$

where $s_1 = (p+k)^2$ and $s_2 = (\bar p + k)^2$. Those for the left-handed lepton ($\sigma = L$) are obtained from the above by the replacement $R \leftrightarrow L$. The tensor wave functions $\epsilon_T^{\mu\nu}$ should satisfy the following conditions for each helicity τ:

$$\epsilon_T^{\mu\nu}(k,\tau) = \epsilon_T^{\nu\mu}(k,\tau)\,, \tag{2.19}$$

$$k_\mu\, \epsilon_T^{\mu\nu}(k,\tau) = k_\nu\, \epsilon_T^{\mu\nu}(k,\tau) = 0\,, \tag{2.20}$$

$$\epsilon_T^{\mu\nu}(k,\tau)\, g_{\mu\nu} = 0\,, \tag{2.21}$$

and they can be normalized as

$$\epsilon_T^{\mu\nu}(k,\lambda)^*\, \epsilon_{T\,\mu\nu}(k,\lambda') = \delta_{\lambda,\lambda'}\,. \tag{2.22}$$

The completeness condition is [12]

$$\sum_\tau \epsilon_T{}^{\alpha\beta}(k,\tau)\, \epsilon_T{}^{\mu\nu}(k,\tau)^* = \frac{1}{2}\left(\kappa^{\mu\alpha}\kappa^{\nu\beta} + \kappa^{\mu\beta}\kappa^{\nu\alpha}\right) - \frac{1}{3}\kappa^{\alpha\beta}\kappa^{\mu\nu}\,, \tag{2.23}$$

where

$$\kappa^{\mu\nu} = -g^{\mu\nu} + \frac{k^\mu k^\nu}{m_T^2} \tag{2.24}$$

for the tensor mass m_T. It should be noted that tensor wave functions are obtained as products of two vector wave functions. The helicity zero ($\tau = 0$) component contains a product of two longitudinary polarized vector wave functions and it behaves as E^4/m_T^4 at high energies. Because of this high-energy behavior, we cannot obtain sensible cross sections for the tensor exchange processes in perturbation theory at energies far above the mass of the tensor boson, while we can regard the interaction (2.17) as an effective one valid at and below the tensor mass scale.

By squaring the amplitude and summing over helicities, we find

$$\sum_{\lambda,\tau,\sigma} |M(\lambda,\tau,\sigma)|^2 = m_Z^2 \left[(g_R f_R^\ell)^2 + (g_L f_L^\ell)^2\right] \Sigma_T(s_1,s_2)\,, \tag{2.25}$$

where we factor out m_Z^2 to account for the mass inverse dimension of the effective couplings f_R^ℓ and f_L^ℓ. The function $\Sigma_T(s_1,s_2)$ is given as

$$\begin{aligned}\Sigma_T(s_1,s_2) = &-\frac{10}{3} - \frac{4s}{m^2} - \frac{s_1+s_2}{s}\left(2 - \frac{10s}{3m^2}\right) + \frac{s_1 s_2}{m^4}\left(\frac{2}{3} + \frac{4m^2}{s}\right) \\ &- \frac{(2s+m^2)(3s+m^2)}{s}\left(\frac{1}{s_1} + \frac{1}{s_2}\right) + \frac{s_1^2+s_2^2}{3\,m^4}\left(1 - \frac{s_1}{s} - \frac{s_2}{s} + \frac{m^2}{s}\right) \\ &- \frac{m^4}{2}\left(\frac{1}{s_1^2} + \frac{1}{s_2^2}\right) + \left(2 + \frac{m^2}{2s}\right)\left\{\frac{s_1}{s_2} + \frac{s_2}{s_1} + \frac{2(s+m^2)^2}{s_1 s_2}\right\}.\end{aligned} \tag{2.26}$$

The Dalitz distribution $\Sigma_T(s_1,s_2)$ is shown in Fig. 4. By comparing with the distributions Figs. 2(a) and 2(b) for the spinless boson case, we find that no subtle cancellation takes place in the Z decay into a tensor ϕ_T and a charged lepton pair.

Upon integration over phase space, we find

$$\Gamma(Z \to \phi_T \ell\bar\ell) = \frac{m_Z^3}{1536\pi^3}\left[(g_R f_R^\ell)^2 + (g_L f_L^\ell)^2\right] F_T(m_T^2/m_Z^2)\,, \tag{2.27}$$

where

$$F_T(x) = \frac{1}{90}(1-x)\left(362\,x^2 + 2897\,x + 2607 - \frac{113}{x} + \frac{7}{x^2}\right) + \frac{1}{3}\ln x\,\left(9\,x^3 + 48\,x^2 + 97\,x + 38\right)$$
$$-(x+1)^2\,(x+4)\left\{-\ln^2 x + \ln^2(1+x) + 2\ln x\,\ln(1+x) + 2\mathrm{Li}_2\,(-x) + 2\mathrm{Li}_2\left(\frac{1}{x+1}\right)\right\}. \quad (2.28)$$

At $x = (59\ \mathrm{GeV}/91.187\ \mathrm{GeV})^2$, we find $F_T(x) = 0.358$. The leptonic width is found as

$$\Gamma(\phi_T \to \ell\bar{\ell}) = \frac{(f_R^\ell)^2 + (f_L^\ell)^2}{320\pi}\, m_T^3\,. \quad (2.29)$$

We hence find

$$\Gamma(\phi_T \to \ell\bar{\ell}) = (3.05\times 10^4\ \mathrm{GeV})\frac{(f_R^\ell)^2 + (f_L^\ell)^2}{(f_R^\ell)^2 + \left(\frac{g_L}{g_R}\right)^2 (f_L^\ell)^2} B(Z \to \phi_T \ell\bar{\ell}) \quad (2.30)$$

for $m_T = 59$ GeV and $\sin^2\theta_W = 0.23$. Since the factor $(g_L/g_R)^2 = [(\frac{1}{2} - \sin^2\theta_W)/\sin^2\theta_W]^2 \approx 1.38$ is almost unity, the required width is insensitive to the specific coupling choice (f_R^ℓ, f_L^ℓ). It is now possible for a narrow resonance of $\sum_{\ell=e,\mu}\Gamma(\phi_T \to \ell\bar{\ell}) \sim 30$ MeV to be produced in Z decays at a branching fraction 10^{-6}.

The factor of ~ 500 difference in the numerical factor for the spinless boson production in Eq. (2.15) and that for the tensor boson production in Eq. (2.30) is semiquantitatively understood as follows. The cancellation in the dominant axial vector contribution $F_A(x)$ to the spinless boson production in Eq. (2.11) gives about a factor of 30 suppression as compared to the vector part $F_V(x)$. The tensor production factor $F_T(x)$ of Eq. (2.27) is even larger than $F_V(x)$ partly because of the spin counting factor $2\times 2+1 = 5$. Finally, the higher dimensionality of the effective tensor coupling in the Lagrangian (2.17) leads to the counting factor of $(m_Z/m_X)^2 \approx 2$. The factors add up to about 300, almost explaining the drastic change in the Z decay branching fractions between the spin-0 and -2 cases.

III. TENSOR EXCHANGE IN e^+e^- COLLIDERS

Effects of a spinless boson exchange in e^+e^- collisions have been rather thoroughly studied in Refs. [6,13]. In this section, we study in detail the consequences of a tensor boson exchange in the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$.

Here we assume that the ϕ_T couplings to the leptons and that with two photons can be parametrized by the Lagrangian

$$\mathcal{L} = \phi_T^{\mu\nu}\left\{\frac{i}{2}\sum_\ell\sum_{\alpha=L,R} f_\alpha^\ell\left[\overline{\psi_{\ell\alpha}}\gamma_\mu(D_\nu\psi_{\ell\alpha}) - (\overline{D_\nu\psi_{\ell\alpha}})\gamma_\mu\psi_{\ell\alpha}\right] + e^2 h\, F_{\mu\lambda}F^\lambda{}_\nu\right\}, \quad (3.1)$$

where the first term in the curly brackets is the $\phi_T\ell\bar{\ell}$ coupling of Eq. (2.17), and $F_{\mu\lambda} = \partial_\mu A_\lambda - \partial_\lambda A_\mu$ is the gauge-invariant field strength of the electromagnetic field. The leptonic and the photonic widths are

$$\Gamma(\phi_T \to \ell\bar{\ell}) = \frac{(f_R^\ell)^2 + (f_L^\ell)^2}{320\pi}\, m_T^3\,, \quad (3.2)$$

$$\Gamma(\phi_T \to \gamma\gamma) = \frac{\pi\alpha^2 h^2}{5}\, m_T^3\,, \quad (3.3)$$

and the total ϕ_T width may be written as

$$\Gamma_T = \sum_\ell \Gamma(\phi_T \to \ell\bar{\ell}) + \Gamma(\phi_T \to \gamma\gamma)\,, \quad (3.4)$$

if other decay modes of ϕ_T are negligible. Summation over lepton flavors should contain $\ell = e$ and μ, but it may or may not include the $\ell = \tau$ case.

As for the ϕ_T propagator, we take the *unitary gauge* form

$$D^{\mu\nu\alpha\beta}(k^2, m_T^2) = \frac{i\sum_\tau \epsilon_T^{\mu\nu}(k,\tau)^*\epsilon_T^{\alpha\beta}(k,\tau)}{k^2 - m_T^2 + i m_T\Gamma_T} \quad (3.5)$$

with the spin summation factor given by Eqs. (2.23) and (2.24). This form behaves as E^2/m_T^4 at high energies ($|k^2| \gg m_T^2$) and cannot be used at energies significantly above the mass shell. It may be regarded as an effective propagator near and below the tensor mass scale, $|k^2| \lesssim m_T^2$.

A. $e^+e^- \to e^+e^-$

We denote the helicity amplitude for the process

$$e^-(p,\sigma) + e^+(\bar{p},\bar{\sigma}) \longrightarrow e^-(k,\lambda) + e^+(\bar{k},\bar{\lambda}) \quad (3.6)$$

as

$$M^{\lambda,\overline{\lambda}}_{\sigma,\overline{\sigma}}, \tag{3.7}$$

where the indices σ, $\overline{\sigma}$, λ, $\overline{\lambda}$ are the helicities in units of $\hbar/2$ [14]. There are six nonvanishing helicity amplitudes

$$M^{-+}_{-+},\ M^{+-}_{-+},\ M^{+-}_{+-},\ M^{-+}_{+-},\ M^{++}_{++},\ M^{--}_{--} \tag{3.8}$$

in the massless electron limit. The most general differential cross section is expressed in terms of these helicity amplitudes as [15]

$$\begin{aligned}\frac{d\sigma}{d\cos\theta\, d\phi} = \frac{1}{256\pi^2 s}\Bigg\{ &(1-P^L_-)\,(1+P^L_+)\left(\left|M^{-+}_{-+}\right|^2+\left|M^{+-}_{-+}\right|^2\right)+(1+P^L_-)\,(1-P^L_+)\left(\left|M^{+-}_{+-}\right|^2+\left|M^{-+}_{+-}\right|^2\right)\\ &+(1+P^L_-)\,(1+P^L_+)\left|M^{++}_{++}\right|^2+(1-P^L_-)\,(1-P^L_+)\left|M^{--}_{--}\right|^2\\ &+2P^T_-P^T_+\cos 2\phi\ \mathrm{Re}\left[\left(M^{-+}_{-+}\right)\left(M^{-+}_{+-}\right)^*+\left(M^{+-}_{-+}\right)\left(M^{+-}_{+-}\right)^*\right]\\ &+2P^T_-P^T_+\sin 2\phi\ \mathrm{Im}\left[\left(M^{-+}_{-+}\right)\left(M^{-+}_{+-}\right)^*+\left(M^{+-}_{-+}\right)\left(M^{+-}_{+-}\right)^*\right]\Bigg\},\end{aligned} \tag{3.9}$$

where θ is the polar angle of the final e^- about the e^- beam axis, ϕ is the azimuthal angle about the e^- beam axis measured from the e^- natural polarization axis in the storage ring, $P^L_\pm$ are the $e^\pm$ beam longitudinal polarizations, and $P^T_\pm$ are the $e^\pm$ beam transverse polarizations. We can set $P^L_+ = P^L_- = 0$ and $P^T_+ = P^T_- = P_T$ at TRISTAN.

The six nonvanishing helicity amplitudes are found as follows :

$$M^{\sigma,-\sigma}_{\sigma,-\sigma} = -s(1+\cos\theta)\left\{\sum_V \left(g^{V\ell\ell}_\sigma\right)^2\left[D_V(s)+D_V(t)\right]+\frac{\left(f^e_\sigma\right)^2 s}{8}\left[(2\cos\theta-1)D_T(s)+\frac{7+\cos\theta}{2}D_T(t)\right]\right\}, \tag{3.10}$$

$$M^{-\sigma,\sigma}_{\sigma,-\sigma} = s(1-\cos\theta)\left\{\sum_V g^{V\ell\ell}_L g^{V\ell\ell}_R D_V(s)+\frac{f^e_L f^e_R\, s}{8}(2\cos\theta+1)D_T(s)\right\}, \tag{3.11}$$

$$M^{\sigma,\sigma}_{\sigma,\sigma} = -2s\left\{\sum_V g^{V\ell\ell}_L g^{V\ell\ell}_R D_V(t)+\frac{f^e_L f^e_R\, s}{16}(5+3\cos\theta)D_T(t)\right\}, \tag{3.12}$$

where $t = -s(1-\cos\theta)/2$, and the propagator factors are

$$D_B(q^2) = \frac{1}{q^2 - m_B^2 + i m_B \Gamma_B \theta(q^2)} \tag{3.13}$$

for $B = V$ (γ or Z) and T. We adopt a chirality index convention $\sigma = -$ for $\sigma = L$, $\sigma = +$ for $\sigma = R$ [14]. The SM couplings are

$$\begin{aligned} g^{\gamma\ell\ell}_L &= g^{\gamma\ell\ell}_R = -e\,,\\ g^{Z\ell\ell}_L &= g_Z\left(-\frac{1}{2}+\sin^2\theta_W\right),\\ g^{Z\ell\ell}_R &= g_Z \sin^2\theta_W\,,\end{aligned} \tag{3.14}$$

with $g_Z = g/\cos\theta_W = e/\sin\theta_W\cos\theta_W$.

The most general differential cross section is easily obtained by inserting the amplitudes (3.10)–(3.12) into (3.9). We show in Fig. 5 the energy dependence of the cross section for the process $e^+e^- \to e^+e^-$ at $\cos\theta = 0$, 0.4 and 0.8, together with the SM predictions. We chose the parameter values $f^\ell_R = 0.0086$, $f^\ell_L = 0$, and $\Gamma_T = 100$ MeV, which give $\sum_{\ell=e,\mu} B(Z \to \phi_T \ell\overline{\ell}) \sim 10^{-6}$ and $B(\phi_T \to \gamma\gamma) \sim 0.7$. At large scattering angles, slightly destructive interference is observed below the resonance peak while above the peak a rather large constructive interference effect is expected. At smaller scattering angles where the t-channel γ exchange amplitude dominates the SM amplitude, the interference below the resonance is found to be constructive. Despite the narrow width of the resonance, its effect is found to be significant in the wide range of the colliding beam energy due to the interference effect. TRISTAN experiments should hence be sensitive to the tensor boson effects even when its width is much narrower than postulated here.

We find that the peak cross section is

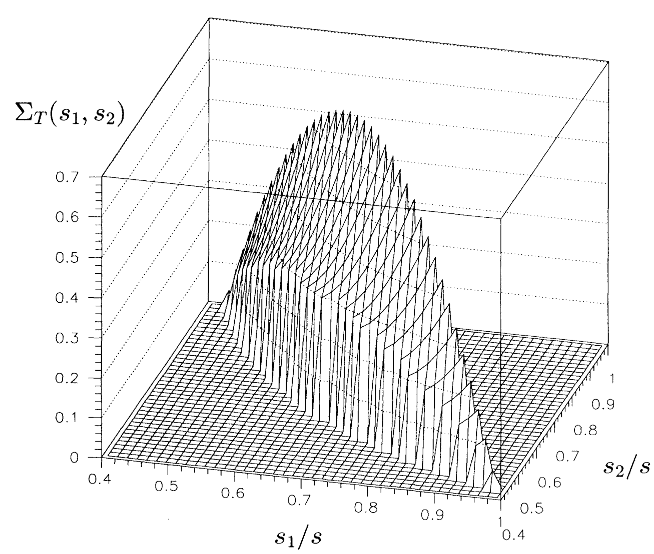


FIG. 4. Dalitz distribution Σ_T of Eq. (2.26) for the decay $Z \to \phi_T \ell_R \overline{\ell_R}$.

$$\sigma(e^+e^- \to \phi_T \to e^+e^-)_{s=m_T^2} = \frac{4\pi}{m_T^2}(2\times 2+1)\frac{\Gamma(\phi_T \to e^+e^-)^2}{\Gamma_T^2} \tag{3.15}$$

$$= \frac{\left[(f_R^e)^2+(f_L^e)^2\right]^2}{5120\pi}\frac{m_T^4}{\Gamma_T^2}\,. \tag{3.16}$$

The polar angle distribution at the peak is

$$\left(\frac{d\sigma}{d\cos\theta}\right)_{s=m_T^2} = \sigma(e^+e^- \to \phi_T \to e^+e^-)_{s=m_T^2}\left\{\frac{5}{8}(1+\cos\theta)^2(1-2\cos\theta)^2 + \frac{5(f_R^e)^2(f_L^e)^2}{\left[(f_R^e)^2+(f_L^e)^2\right]^2}\cos\theta(1-2\cos^2\theta)\right\}. \tag{3.17}$$

B. $e^+e^- \to \mu^+\mu^-$, $\tau^+\tau^-$

For the process

$$e^-(p,\sigma)+e^+(\overline{p},\overline{\sigma}) \longrightarrow \ell\,(k,\lambda)+\overline{\ell}\,(\overline{k},\overline{\lambda}) \tag{3.18}$$

with $\ell=\mu$ or τ, only the first four helicity amplitudes of (3.8) are nonvanishing in the massless lepton limit. They are obtained from Eqs. (3.10) and (3.11) simply by replacing the leptonic coupling factors and by dropping the t-channel exchange contribution:

$$M_{\sigma,-\sigma}^{\sigma,-\sigma} = -s\,(1+\cos\theta)\left\{\sum_V g_\sigma^{Vee}\,g_\sigma^{V\ell\ell}\,D_V(s)+\frac{1}{8}\,f_\sigma^e f_\sigma^\ell s\,(2\cos\theta-1)\,D_T(s)\right\}, \tag{3.19}$$

$$M_{\sigma,-\sigma}^{-\sigma,\sigma} = s\,(1-\cos\theta)\left\{\sum_V g_\sigma^{Vee}\,g_{-\sigma}^{V\ell\ell}\,D_V(s)+\frac{1}{8}\,f_\sigma^e\,f_{-\sigma}^\ell s(2\cos\theta+1)\,D_T(s)\right\}. \tag{3.20}$$

The most general differential cross section is then obtained by inserting the above helicity amplitudes into the formula (3.9). The peak cross section is

$$\sigma(e^+e^- \to \phi_T \to \ell^+\ell^-)_{s=m_T^2} = \frac{4\pi}{m_T^2}(2\times 2+1)\frac{\Gamma(\phi_T \to e^+e^-)\,\Gamma(\phi_T \to \ell^+\ell^-)}{\Gamma_T^2}, \tag{3.21}$$

and the polar angle distribution at the peak is

$$\left(\frac{d\sigma}{d\cos\theta}\right)_{s=m_T^2} = \sigma(e^+e^- \to \phi_T \to \ell^+\ell^-)_{s=m_T^2} \times\left\{\frac{5}{8}(1+\cos\theta)^2(1-2\cos\theta)^2+\frac{5\,f_R^e f_L^e f_R^\ell f_L^\ell}{\left[(f_R^e)^2+(f_L^e)^2\right]\left[(f_R^\ell)^2+(f_L^\ell)^2\right]}\cos\theta\,(1-2\cos^2\theta)\right\}. \tag{3.22}$$

We show in Fig. 6 the energy dependence of the differential cross section for the process $e^+e^- \to \mu^+\mu^-$ at $\cos\theta=0$, together with the SM prediction depicted by the dashed line. The tensor couplings to left-handed leptons have been set to zero, $f_L^e=f_L^\mu=0$, as before. A significant interference effect is again expected for the postulated resonance parameters [$m_T=59$ GeV and $\Gamma(\phi_T \to e^+e^-)=15$ MeV]. The interference below the resonance peak is constructive (destructive) if the tensor couplings to electrons (f_R^e) and those to muons (f_R^μ) have the same (opposite) sign. It is clear from the figure that TRISTAN experiments should be sensitive to a tensor resonance with much smaller leptonic width.

C. $e^+e^- \to \gamma\gamma$

We denote the helicity amplitude for the process

$$e^-(p,\sigma)+e^+(\overline{p},\overline{\sigma}) \longrightarrow \gamma\,(k_1,\lambda_1)+\gamma\,(k_2,\lambda_2) \tag{3.23}$$

as

$$M_{\sigma\,,\,\overline{\sigma}}^{\lambda_1,\lambda_2}\,, \tag{3.24}$$

where $\lambda_i=\pm$ denotes the γ helicities in units of $\hbar$. We find in the convention of Ref. [14] that only the following four helicity amplitudes ($\sigma=\pm$ and $\lambda=\pm$),

$$M_{\sigma,-\sigma}^{\lambda,-\lambda} = 2\,\sigma\lambda\,e^2\frac{1+\sigma\lambda\cos\theta}{\sin\theta}\left[1+\frac{f_\sigma^e h\,s^2\sin^2\theta}{8\,(m_T^2-s-im_T\Gamma_T)}\right], \tag{3.25}$$

are nonvanishing in the massless electron limit. It is worth noting that both the SM and the tensor boson exchange contributions satisfy the *helicity conservation* condition, $\lambda_2 = -\lambda_1$, for the two outgoing photons. This leads to significant interference effects between the two contributions.

The general differential cross section is again obtained by inserting the helicity amplitudes (3.25) into the generic formula (3.9). For $P_+^L = P_-^L = 0$ and $P_+^T = P_-^T = P_T$, we find, more explicitly,

$$\begin{aligned}\frac{d\sigma}{d\cos\theta\, d\phi} = \frac{\alpha^2}{s}\Bigg\{ & \frac{1+\cos^2\theta}{\sin^2\theta} + \frac{(f_R^e + f_L^e)\, hs^2(m_T^2 - s)}{8\left[(m_T^2-s)^2 + (m_T\Gamma_T)^2\right]}(1+\cos^2\theta) \\ & + \frac{[(f_R^e)^2 + (f_L^e)^2]h^2 s^4}{128\left[(m_T^2-s)^2 + (m_T\Gamma_T)^2\right]}\sin^2\theta\,(1+\cos^2\theta) \\ & - P_T^2\left[1 + \frac{(f_R^e + f_L^e)\, hs^2(m_T^2-s)}{8\left[(m_T^2-s)^2+(m_T\Gamma_T)^2\right]}\sin^2\theta + \frac{f_R^e f_L^e h^2 s^4}{64\left[(m_T^2-s)^2+(m_T\Gamma_T)^2\right]}\sin^4\theta\right]\cos 2\phi \\ & - P_T^2\frac{(-f_R^e + f_L^e)\, hs^2 m_T\Gamma_T}{8\left[(m_T^2-s)^2+(m_T\Gamma_T)^2\right]}\sin^2\theta\,\sin 2\phi \Bigg\}. \end{aligned} \tag{3.26}$$

The peak cross section is

$$\sigma(e^+e^- \to \gamma\gamma)_{s=m_T^2} = \frac{4\pi}{m_T^2}(2\times 2+1)\frac{\Gamma(\phi_T \to e^+e^-)\,\Gamma(\phi_T\to\gamma\gamma)}{\Gamma_T^2}, \tag{3.27}$$

and the angular distribution at the peak is

$$\begin{aligned}\left(\frac{d\sigma}{d\cos\theta\, d\phi}\right)_{s=m_T^2} &= \sigma(e^+e^-\to\gamma\gamma)_{s=m_T^2}\frac{5}{8\pi}\sin^2\theta \\ &\times\left\{1+\cos^2\theta + P_T\left[\frac{2f_R^e f_L^e\sin^2\theta}{(f_R^e)^2+(f_L^e)^2}\cos 2\phi + \frac{16(-f_R^e+f_L^e)\Gamma_T}{[(f_R^e)^2+(f_L^e)^2]\, hm_T^3}\sin 2\phi\right]\right\}. \end{aligned} \tag{3.28}$$

Shown in Fig. 7 is the energy dependence of the expected differential cross section at $\cos\theta = 0$ for $m_T = 59$ GeV, $\Gamma(\phi_T \to e^+e^-) = 15$ MeV, and $\Gamma(\phi_T \to \gamma\gamma) = 70$ MeV ($\Gamma_T = 100$ MeV). The SM prediction is given by the dashed line. A huge interference effect is expected for the above resonance parameters. The effect below the resonance is constructive if the couplings f_R^e and h have the common sign, whereas it is destructive otherwise. We may conclude from Figs. 5, 6, and 7 that the tensor resonance effect can be most clearly observed in the $e^+e^- \to \gamma\gamma$ channel.

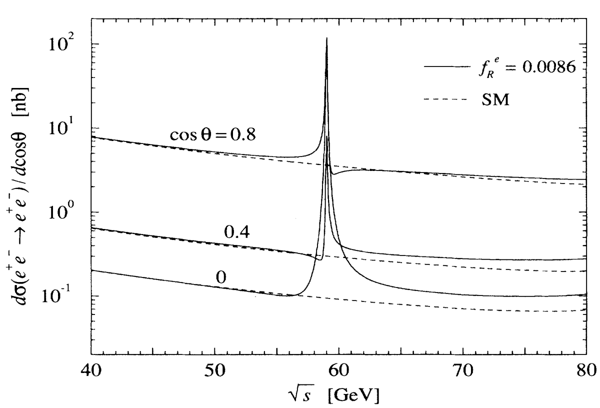


FIG. 5. The energy dependence of the differential cross section $d\sigma(e^+e^- \to e^+e^-)/d\cos\theta$ at $\cos\theta = 0$, 0.4, and 0.8 for $f_R^e = 0.0086$, $f_L^e = 0$, and $\Gamma_T = 100$ MeV. The dashed lines give the SM predictions for $e^2 = 4\pi\overline{\alpha}(m_Z^2)$ with $\overline{\alpha}(m_Z^2) = 1/128$, $\sin^2\theta_W = 0.23$, $m_Z = 91.187$ GeV, and $\Gamma_Z = 2.5$ GeV.

IV. LARGE TENSOR MASS LIMITS

We found in the previous section that the effects of a tensor boson exchange can be observed at energies far away from the resonance peak due to the interference be-

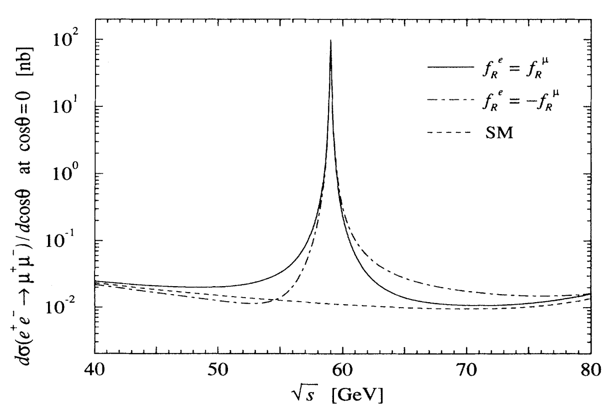


FIG. 6. The energy dependence of the differential cross section $d\sigma(e^+e^- \to \mu^+\mu^-)/d\cos\theta$ at $\cos\theta = 0$ for $f_R^e = \pm f_R^\ell = 0.0086$ and $f_L^e = f_L^\ell = 0$. The dashed line gives the SM prediction for the SM parameters of Fig. 5.

tween the tensor exchange amplitudes and the SM ones. This is in contrast with the spinless boson exchange case [6] where a significant interference effect is expected only in the $e^+e^- \to e^+e^-$ channel. The difference arises because the tensor boson coupling to leptons can be helicity conserving just like the SM gauge boson couplings while the spinless boson coupling should necessarily flip the lepton helicities. Therefore, e^+e^- collider experiments can have sensitivity to tensor boson exchange effects even when its mass is far above the colliding beam energy.

In the large m_T limit ($s/m_T^2 \ll 1$), the ϕ_T exchange contributions to the processes $e^+e^- \to \ell\bar{\ell}$ are related to those of the standard contact $ee\ell\ell$ interaction [7] at fixed energy and scattering angle. Note that we cannot simply substitute the limits on the combination $f_\sigma^e f_\sigma^\ell/m_T^2$ by those of the corresponding contact terms independent of energies and angles, because the ϕ_T exchange gives effective four-fermion operators of spin 2 and dimension 8 whereas the standard contact terms have spin 1 and dimension 6.

On the other hand, the dimension-6 contact $ee\gamma\gamma$ interaction terms as listed by Ref. [16] give vanishing matrix elements for the process $e^+e^- \to \gamma\gamma$. It has been found [8] that the lowest dimensional local operator that is chiral invariant and has nonvanishing $e^+e^- \to \gamma\gamma$ matrix elements has dimension 8 and spin 2. Therefore, in the large m_T limit, the ϕ_T exchange contribution reduces to the dimension-8 contact $ee\gamma\gamma$ term in the process $e^+e^- \to \gamma\gamma$.

Let us recall the standard dimension-6 contact interaction for leptons [7,17]:

$$\mathcal{L} = \frac{g^2}{\Lambda^2} \sum_{(\ell,\ell')} \sum_{\alpha=L,R} \sum_{\beta=L,R} \eta_{\alpha\beta}^{\ell\ell'} S^{\ell\ell'} \overline{\psi_{\ell\alpha}} \gamma^\mu \psi_{\ell\alpha} \overline{\psi'_{\ell\beta}} \gamma_\mu \psi'_{\ell\beta}, \tag{4.1}$$

where $g^2 = 4\pi$ by convention, $|\eta_{\alpha\beta}^{\ell\ell'}| = 1$ or 0 for special cases, and

$$S^{\ell\ell'} = \begin{cases} 1 & \text{for } \ell \neq \ell' \\ \frac{1}{2} & \text{for } \ell = \ell' \end{cases} \tag{4.2}$$

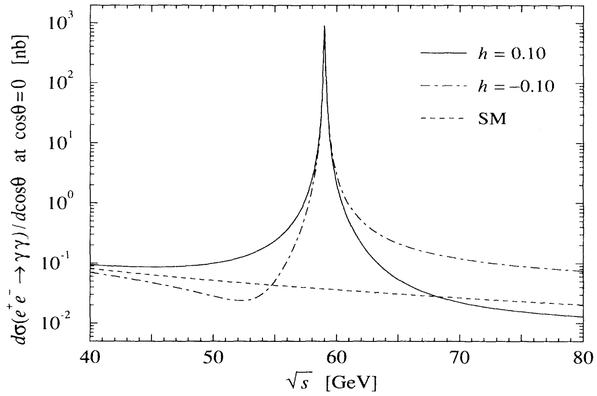


FIG. 7. The energy dependence of the differential cross section $d\sigma(e^+e^- \to \gamma\gamma)/d\cos\theta$ at $\cos\theta = 0$ for $f_R^e = 0.0086$, $f_L^e = 0$, and $h = \pm 0.10$. The dashed line gives the SM prediction for $e^2 = 4\pi\alpha$ with $\alpha = 1/137$.

is the statistical factor.

We find that the helicity amplitudes for the process $e^+e^- \to \ell\bar{\ell}$ (3.19) and (3.20) reproduce those with the above contact interaction terms by the substitutions

$$\frac{f_\sigma^e f_\sigma^\ell s}{8}(2\cos\theta - 1)D_T(s) \longrightarrow \frac{g^2\eta_{\sigma\sigma}^{e\ell}}{\Lambda^2}, \tag{4.3}$$

$$\frac{f_\sigma^e f_{-\sigma}^\ell s}{8}(2\cos\theta + 1)D_T(s) \longrightarrow \frac{g^2\eta_{\sigma,-\sigma}^{e\ell}}{\Lambda^2}. \tag{4.4}$$

In the large mass limit, the following relationships hold:

$$\frac{f_\sigma^e f_\sigma^\ell s}{8m_T^2}(1 - 2\cos\theta) = \frac{g^2\eta_{\sigma\sigma}^{e\ell}}{\Lambda^2}, \tag{4.5}$$

$$-\frac{f_\sigma^e f_{-\sigma}^\ell s}{8m_T^2}(1 + 2\cos\theta) = \frac{g^2\eta_{\sigma,-\sigma}^{e\ell}}{\Lambda^2}. \tag{4.6}$$

The helicity amplitudes for the process $e^+e^- \to e^+e^-$ (3.10)–(3.12) reduce to those with the contact $eeee$ term via the substitutions

$$\frac{(f_\sigma^e)^2 s}{8}\left[(2\cos\theta - 1)D_T(s) + \frac{1}{2}(7 + \cos\theta)D_T(t)\right] \longrightarrow \frac{2g^2\eta_{\sigma,\sigma}^{ee}}{\Lambda^2}, \tag{4.7}$$

$$\frac{f_L^e f_R^e s}{8}(2\cos\theta + 1)D_T(s) \longrightarrow \frac{g^2\eta_{LR}^{ee}}{\Lambda^2}, \tag{4.8}$$

$$\frac{f_L^e f_R^e s}{16}(5 + 3\cos\theta)D_T(t) \longrightarrow \frac{g^2\eta_{LR}^{ee}}{\Lambda^2}. \tag{4.9}$$

In the $m_T^2 \gg s$ limit, the substitution rule (4.7) gives the identification

$$-\frac{5(f_\sigma^e)^2 s}{16m_T^2}(1 + \cos\theta) = \frac{2g^2\eta_{\sigma\sigma}^{ee}}{\Lambda^2}. \tag{4.10}$$

However, the rules (4.8) and (4.9) do not allow us to relate the parameters η_{LR}^{ee} consistently with the corresponding terms of the ϕ_T exchange parameters even in the large m_T limit at fixed energy and scattering angle.

The lowest dimension chirality-conserving $ee\gamma\gamma$ operator which gives nonvanishing $e^+e^- \to \gamma\gamma$ matrix elements is found to be [8]

$$\mathcal{L} = \frac{2ie^2}{\Lambda^4} F^{\mu\sigma} F_\sigma{}^\nu \sum_{\alpha=L,R} \eta_\alpha \overline{\psi}_\alpha \gamma_\mu \partial_\nu \psi_\alpha. \tag{4.11}$$

The helicity amplitudes with the above contact terms are

$$M_{\sigma,-\sigma}^{\lambda,-\lambda} = 2\sigma\lambda e^2 \frac{1+\sigma\lambda\cos\theta}{\sin\theta}\left[1+\frac{\eta_\sigma s^2 \sin^2\theta}{4\Lambda^4}\right]. \tag{4.12}$$

By comparing the above amplitudes and the ϕ_T exchange ones (3.25), we find the substitution rule

$$\frac{f_\sigma^e h}{8(m_T^2 - i m_T \Gamma_T - s)} \longrightarrow \frac{\eta_\sigma}{4\Lambda^4}. \tag{4.13}$$

In the large m_T limit ($m_T^2 \gg s$), we can make the identification

$$\frac{f_\sigma^e h}{8 m_T^2} = \frac{\eta_\sigma}{4\Lambda^4}, \tag{4.14}$$

which is valid at all energies and angles.

The relationships between the tensor exchange amplitudes and the contact $ee\ell\ell$ and $ee\gamma\gamma$ interactions may be useful to estimate bounds on the ϕ_T parameters in the $m_T^2 \gg s$ limit from the existing limit on the contact interactions.

V. CONCLUSIONS

Assuming that high mass $\gamma\gamma$ events observed at LEP are caused by a $\gamma\gamma$ resonance which is produced via its leptonic couplings, we explored its consequences at TRISTAN. We examined both spin-0 and spin-2 cases as the simplest possibilities.

We showed that, in order to account for the branching fraction $B(Z \to \phi_{S,P}\ell\bar{\ell})$ of the order of 10^{-6}, the couplings of a spin-0 boson to leptons have to be so large that the boson width should be comparable to its mass of about 60 GeV. This is because of a strong cancellation in the decay amplitudes for the almost axial vector $Z\ell\bar{\ell}$ coupling of the SM, which causes a destructive interference between the two Feynman diagrams of Fig. 1. As a result, the spin-0 case is excluded whether it is a scalar (ϕ_S), or a pseudoscalar (ϕ_P), or their doublet.

For the case of a spin-2 boson (ϕ_T), no such cancellation arises because it has chirality-conserving couplings to a lepton. We found that a narrow resonance of about 30 MeV leptonic width can be produced in Z decays at the required branching fraction level. A higher spin resonance may improve the situation further by means of spin and dimensional counting.

If there is such a $\gamma\gamma$ tensor resonance around 60 GeV, we can observe its effect at TRISTAN. We gave complete helicity amplitudes including tensor boson exchange for the processes $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$. Differential cross sections for arbitrarily polarized beams are easily obtained from the amplitudes. The total cross sections and angular distributions behave quite differently from the spin-0 resonance cases [6]. The resonance spin can hence be determined from these distributions.

We further noted that the tensor resonance effects can be observed at TRISTAN even if the colliding beam energy is far away from the resonance peak, due to the interference effects between the tensor exchange amplitudes and the SM ones. The experiments should hence be sensitive to the effects of a tensor resonance of a much smaller leptonic width or a much larger mass than has been postulated as a possible origin of the 60 GeV $\gamma\gamma$ events observed at LEP.

Finally, we found that the tensor boson exchange effects in the processes $e^+e^- \to \ell^+\ell^- (\ell \neq e)$ and $\gamma\gamma$ reduce to those of the contact higher dimensional interactions in the high tensor boson mass limit. We identified the substitution rules between the tensor exchange effects and the effects of dimension-6 contact $ee\ell\ell$ couplings [7] for the process $e^+e^- \to \ell^+\ell^-$, and those between the tensor exchange and the dimension-8 contact $ee\gamma\gamma$ interactions [8] for the process $e^+e^- \to \gamma\gamma$.

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