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Four-fermion processes at future e^+e^- colliders as a probe of new resonant structures

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Possible oblique effects from vector particles that are strongly coupled to the known gauge bosons are calculated for the case of final hadronic states produced at future e^+e^- colliders, using a formalism that was recently proposed and that exploits the information and the constraints provided by CERN LEP 1 results. Combining the hadronic channels with the previously analyzed leptonic ones we derive improved limits for the masses of the resonances that, in technicolorlike cases, would range from one to two TeV for a 500 GeV linear collider, depending on the assumed theoretical constraints.

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The possibility of using high precision data from the CERN e^+e^- collider LEP 1 to derive information or to set stringent bounds on technicolor models, has been thoroughly investigated in recent times, following the original proposal of Peskin and Takeuchi [1]. As is known, the relevant effect is a virtual one-loop contribution of the so-called [2] oblique type to the quantity defined as S in Ref. [1].

Technically speaking, the calculation of S is made easier by the fact that the combination of spectral functions involved has a rather exceptional asymptotic convergence, being the difference of a vector and an axial vector term, and this allows the use of simple dispersion relations, i.e., without unknown extra subtraction constants. This nice feature would not be present in general in different kinematical configurations, e.g., away from the Z resonance, for other oblique corrections of similar type, and an analogous calculation of technicolorlike effects would require some extra ingredient or *ad hoc* assumptions that might bias the theoretical outcome.

In a recent publication [3] we actually proposed a general formalism to calculate the relevant oblique contributions to a number of processes in future higher energies e^+e^- experiments. The main idea was that of expressing the various effects in the form of a once-subtracted dispersion integral, and of fixing the necessary subtraction constants by suitable model-independent LEP 1 results. In this way, we were led to a compact "representation" of several observables. In particular, we concentrated our preliminary analysis on the case of final leptonic states and more precisely on the three quantities: (a) the cross section for muon production at c.m. energy $\sqrt{q^2}$, $\sigma_\mu(q^2)$; (b) the related forward-backward asymmetry $A_{\mathrm{FB},\mu}(q^2)$; (c) the (conventionally defined) final τ polarization asymmetry $A_\tau(q^2)$ or, equivalently, the longitudinal polarization asymmetry for final *lepton* production $A_{LR,l}(q^2)$ whose theoretical expressions coincide in our scheme.

Starting from the tree-level expressions of (a), (b), (c) and making use of the by now conventional formalism based on the introduction of the two parameters $\epsilon_{1,3}$ that allow us to interpret LEP 1 leptonic data in a model-independent way [4], we were able to write for the oblique (SE=self-energy) corrections the following approximate formulas, valid at the one loop level:

$$\sigma_\mu^{\rm SE}(q^2)=\frac{4\pi q^2}{3}\left\{\left[\frac{\alpha(M_Z^2)}{q^2}\right]^2[1+2D_\gamma(q^2)]+\frac{1}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}\left[\frac{3\Gamma_l}{M_Z}\right]^2\left[1-2D_Z(q^2)-\frac{16s_1^2v_1}{1-v_1^2}D_{\gamma Z(q^2)}\right]\right\}, \quad (1)$$

$$A_{{\rm FB},\mu}^{\rm SE}(q^2)=\frac{3}{4}\left[\frac{3q^2\sigma_\mu(q^2)}{4\pi}\right]^{-1}\left\{6\alpha(M_Z^2)\frac{\Gamma_l}{M_Z}\frac{q^2(q^2-M_Z^2)}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}\{1+D_\gamma(q^2)-D_Z(q^2)]\}\right\}, \quad (2)$$

$$A_\tau^{\rm (SE)}(q^2)\equiv A_{LR,l}^{\rm (SE)}=\left[\frac{3q^2\sigma_\mu(q^2)}{4\pi}\right]^{-1}A(M_Z^2)\left\{\left[6\alpha(M_Z^2)\frac{\Gamma_l}{M_Z}\frac{q^2(q^2-M_Z^2)}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}\right.\right.$$
$$\left.\left.+18\left[\frac{\Gamma_l}{M_Z}\right]^2\frac{q^4}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}\right]\left[1-\frac{8s_1^2}{A(M_Z^2)}D_{\gamma Z}(q^2)\right]\right\}. \quad (3)$$

Here Γ_l is the leptonic Z width, $\alpha(M_Z^2)=[1\pm0.001]/128.87$ [5], $A(M_Z^2)$ is defined as

$$A(M_Z^2)\equiv\frac{2[1-4s_{\rm eff}^2(M_Z^2)]}{1+[1-4s_{\rm eff}^2(M_Z^2)]^2} \quad (4)$$

with $s_{\rm eff}^2(M_Z^2)$ measured by the various asymmetries at LEP 1 and the SLAC Linear Collider (SLC), and

$$D_\gamma(q^2)\equiv\Delta\alpha(q^2)-\Delta\alpha(M_Z^2)=-\frac{q^2-M_Z^2}{\pi}\mathcal{P}\int_0^\infty\frac{ds\,{\rm Im}F_\gamma(s)}{(s-q^2)(s-M_Z^2)}, \quad (5)$$

$$D_Z(q^2)\equiv{\rm Re}[I_Z(q^2)-I_Z(M_Z^2)]=\frac{q^2-M_Z^2}{\pi}\mathcal{P}\int_0^\infty\frac{ds\,s\,{\rm Im}F_{ZZ}(s)}{(s-q^2)(s-M_Z^2)^2}, \quad (6)$$

$$D_{\gamma Z}(q^2)\equiv{\rm Re}[\Delta\bar\kappa'(q^2)-\Delta\bar\kappa'(M_Z^2)]=\frac{q^2-M_Z^2}{\pi}\mathcal{P}\int_0^\infty\frac{ds\,{\rm Im}F_{\kappa'}(s)}{(s-q^2)(s-M_Z^2)}, \quad (7)$$

$$(F'_\kappa=c_1/s_1F_{Z\gamma}, s_1^2c_1^2=\pi\alpha/\sqrt{2}G_\mu M_Z^2,\ \ s_1^2=1-c_1^2\simeq0.217,\ \ v_1=1-4s_1^2).$$

Equations (1)–(3) provide a representation of the leptonic observables of e^+e^- annihilation where the full effect of the oblique corrections is made explicit in the form of a subtracted dispersion relation, thus calculable for models of both perturbative and of nonperturbative type, with the subtraction constants provided by model-independent LEP 1 data. Note that, to obtain properly gauge-invariant expressions, one still needs to add the correct amount of extra vertices and boxes [6], as discussed in Ref. [3], to compensate for the intrinsically not gauge-invariant nature of the transverse self-energies, defined following the convention

$$A_{ij}(q^2)\equiv A_{ij}(0)+q^2F_{ij}(q^2),\ \ i,j=\gamma,Z\ . \quad (8)$$

Starting from Eqs. (1)–(3) and (5)–(7) we calculated in Ref. [3] the possible effects of a couple of vector (V) and axial vector (A) resonances with masses larger than $\sqrt{q^2}$, strongly coupled to the photon and to the Z. We assumed a "technicolorlike" framework but only exploited the validity of the *second* Weinberg sum rule [7]. We did not use the model-dependent information provided by the first Weinberg sum rule. However, we retained one very general consequence of it, i.e., the positivity of S, which was ensured by the choice $M_A>M_V$. Taking into account the LEP 1 constraint [8] on the S parameter, we derived observability limits for $M_{V,A}$ in the TeV range for a realistic e^+e^- linear collider of 500 GeV c.m. energy [9]. This was an encouraging preliminary result, particularly since only the final leptonic channels were fully exploited.

This short paper has two purposes. The first one is that of enlarging the previous study by including the potentially copious information provided by the analysis of final hadronic states. The second one is that of emphasizing the relevance of some special theoretical assumptions to fix the derived mass limits, in particular, of investigating the consequences of relaxing completely the two Weinberg sum rules, while still retaining the experimental constraint provided by the LEP 1 limits on the S parameter.

The investigation of the hadronic channels can be easily performed following the prescriptions of Ref. [3]. We shall briefly sketch here the derivation of the relevant formulas for the "basic" cases of the two cross sections for production of u-type and d-type quarks, $\sigma_{u,d}(q^2)$. With this purpose, we start from the expressions of these quantities at the tree level:

$$\sigma^{(0)}_{u,d}(q^2)=N^{(0)}_{u,d}\left[\frac{4}{3}\pi q^2\right]\left\{\left[\frac{Q_{u,d}\alpha_0}{q^2}\right]^2+\left[\frac{G^0_\mu\sqrt{2}M^2_{0Z}}{16\pi}\right]^2\frac{16[(g^0_{V,u,d})^2+(g^0_{A,u,d})][(g^0_{V,l})^2+(g^0_{A,l})]}{D^2_{0Z}}\right.$$
$$\left.-2Q_{u,d}\frac{\alpha_0 G^0_\mu\sqrt{2}M^2_{0Z}}{16\pi q^2}4g^0_{V,l}g^0_{V,u,d}\mathrm{Re}\frac{1}{D_{0Z}}\right\}, \tag{9}$$

where $N_{u,d}$ is the color factor, $g_{V,A,f}$ are conventionally defined, i.e., $g_{A_0,f}=T_{3L,f}$ and $g_{V0,f}=T_{3L,f}-2Q_f s_0^2$, $G_{\mu 0}$ is the (bare) Fermi muon decay coupling and $D_{0Z}=q^2-M^2_{0Z}$ [the tree level equality $\alpha_0/s_0^2c_0^2=(\sqrt{2}/\pi)G_{\mu 0}M^2_{0Z}$ has been used].

When moving to one loop, one has to redefine the Fermi coupling, the QED coupling, the bare mass M_Z, the photon and Z propagators and the various fermion couplings $g_{V,A,f}$. Then, vertex corrections and boxes should be correctly included. For the specific purposes of this paper, that is, only dealing with oblique corrections, these terms will not be explicitly calculated. Thus, in the redefinition of the Fermi coupling, only the oblique content $A_{WW}(0)/M_W^2$ will be retained. Analogously, for the vector couplings we shall stick to the notation of a previous paper [10] and write, following essentially the Kennedy and Lynn approach, [11]

$$\frac{g_{V,f}}{g_{A,f}}=1-4|Q_f|^2 s_f^2(q^2)\quad\text{with } s_f^2(q^2)=s_1^2[1+\Delta\bar{\kappa}'_f(q^2)]\ . \tag{10}$$

The quantity $\Delta\bar{\kappa}'_f(q^2)$ can be decomposed into a universal self-energy component $\Delta\bar{\kappa}'$ and a (light) fermion-dependent vertex correction, i.e. (omitting boxes),

$$\Delta\bar{\kappa}'_f(q^2)=\Delta\bar{\kappa}'(q^2)+\delta'_f \tag{11}$$

with δ'_f (to be from now on neglected) defined in Ref. [10] and $\Delta\bar{\kappa}'(q^2)$ fixed by the convention

$$S^2_{\mathrm{EFF}}(M_Z^2)=s_1^2[1+\Delta\bar{\kappa}'(M_Z^2)+\delta'_l]\ . \tag{12}$$

The procedure for deriving compact expressions for the various self-energy contributions at one loop now follows essentially the same lines as in the case of Ref [3]. In fact, the pure photon contribution will generate the usual term $\simeq D_\gamma(q^2)$. From the Z contribution, using the definition [4]

$$\Gamma_l=\frac{G_\mu M_Z^3}{24\pi\sqrt{2}}[1+\epsilon_1][1+\{1-4s^2_{\mathrm{EFF}}(M_Z^2)\}^2] \tag{13}$$

with

$$\epsilon_1\equiv-\frac{A_{WW}(0)}{M_W^2}+\frac{A_{ZZ}(0)}{M_Z^2}+\text{vertices}\cdots \tag{14}$$

as expression containing Γ_l, $D_Z(q^2)$, $D_{\gamma Z}(q^2)$, and $s^2_{\mathrm{EFF}}(M_Z^2)$ will be originated. Finally, from the $\gamma-Z$ interference, a combination of the previous pure photon and pure Z case parameters will appear. In practice, the main difference between the self-energy content of $\sigma_{u,d}$ and that of σ_μ will come from the relative weights of the various D_γ, D_Z, $D_{\gamma Z}$ contributions, due to the various electric charges Q_f that enter both as coefficients of α and as coefficients of g_V/g_A in Eq. (9).

With these premises, it becomes relatively simple to derive the explicit expressions of the desired one-loop contributions to $\sigma_{u,d}$. Neglecting systematically numerically irrelevant contributions, one obtains the following simple formulas ($\sigma_f^{(1)}$ denotes the quantity at one loop):

$$\sigma^{(1)}_{u,d}=N^{(1)}_u\left[\frac{4}{3}\pi q^2\right]\left\{Q^2_{u,d}\left[\frac{\alpha(M_Z^2)}{q^2}\right]^2[1+2D_\gamma(q^2)]\right.$$
$$+\frac{1}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}\left[\frac{3\Gamma_l}{M_Z}\right]^2[1+v^2_{u,d1}][1-2D_Z(q^2)]-2Q_{u,d}\alpha(M_Z^2)$$
$$\left.\times\frac{3\Gamma_l}{M_Z}\frac{q^2-M_Z^2}{q^2}\frac{1}{(q^2-M_Z^2)^2+M_Z^2\Gamma_Z^2}v_{u,d,1}v_1\left[1-\left[\frac{4s_1^2}{v_1}+4|Qu,d|\frac{s_1^2}{v_{u,d1}}\right]D_{\gamma Z}(q^2)\right]\right\} \tag{15}$$

where $N_f^{(1)}$ is the color QCD corrected factor and we used the generalized notation

$$v_{f1}\equiv1-4|Q_f|s_1^2,\quad f=u,d \tag{16}$$

(and $v_{l1}\equiv v_1$). Starting from the "basic" quantities, Eqs. (1) and (15), it is now straightforward to derive the corresponding expressions of a certain number of hadronic observables. We have considered here the theoretical expressions of the following "candidates" to reveal potential self-energy effects. (I) $R^{(5)}(q^2)$, the ratio $\sigma^{(5)(q^2)}/\sigma_\mu(q^2)$ between the cross sections for production of the five lighter (u,d,s,c,b) quarks and for muon production; (II) $A_{LR}^{(5)}(q^2)$, the longitudinal polarization asymmetry for final hadronic states of the previous type; (III) $R_{b,\mu}(q^2)$, the ratio $\sigma_b(q^2)/\sigma_\mu(q^2)$ between b quark and muon production; (IV) $A_{\mathrm{FB},b}(q^2)$, the forward-backward asymmetry for b-quark production. In addition to the previous "old-fashioned" quantities, we have also calculated, assuming a (copious) top production at $\sqrt{q^2}=500$ GeV, the theoretical expression of a number of related observables. In particular, we have considered here (V)

$R^{(6)}(q^2)$, $A_{LR}^{(6)}(q^2)$, and $R_{t,\mu}(q^2)$, defined in analogy with (I), (II), and (III), and (VI) $R_{b,t}^{(5),(6)}(q^2)$ [the ratios $\sigma_{b,t}(q^2)/\sigma^{(5),(6)}(q^2)$] and $A_{FB,t}(q^2)$.

For all the previous observables from (I) to (VI), it is not difficult to write the expressions at one loop that generalize those of Ref. [3]. But, in the actual process of doing that, one easily realizes that *a priori* not all cases seem equally promising. In particular, assuming "realistic" experimental accuracies (i.e., of the kind discussed in previous analyses [9]) for the various cross sections and their ratios, it turns out that the weights of the various $D_\gamma, D_Z, D_{\gamma,Z}$ contributions (that are rather different in the various observables) are systematically "small" in the cases (IV)–(VI), leading to practically unobservable effects. For this reason, we concentrated our attention on the quantities (I)–(III) only. Ignoring as usual several irrelevant terms, we were led in conclusion to the following set of expressions that include the full effect of the oblique corrections at one loop:

$$R^{(5)(\mathrm{SE})}(q^2)=a_0[1+a_\gamma D_\gamma+a_Z D_Z+a_{\gamma Z}D_{\gamma Z}]\ , \qquad (17)$$

$$R_{b,\mu}^{(\mathrm{SE})}(q^2)=b_0[1+b_\gamma D_\gamma+b_Z D_Z+b_{\gamma Z}D_{\gamma Z}]\ , \qquad (18)$$

$$A_{LR}^{(5)(\mathrm{SE})}(q^2)=c_0[1+c_\gamma D_\gamma+c_Z D_Z+c_{\gamma Z}D_{\gamma Z}]\ , \qquad (19)$$

where the analytic expressions of the various coefficients can be derived in a straightforward way and their numerical values for $\sqrt{q^2}=500(190)$ GeV are given as

$$a_0=5.59(6.84),\quad a_\gamma=-0.61(-0.76),\quad a_Z=-0.84(-1.12),\quad a_{\gamma Z}=-0.26(-0.32)\ ,$$
$$b_0=0.88(1.16),\quad b_\gamma=-1.10(-1.17),\quad b_Z=-1.21(-1.41),\quad b_{\gamma Z}=-0.80(-0.78)\ ,$$
$$c_0=0.61,\quad c_\gamma=-0.42,\quad c_Z=-0.27,\quad c_{\gamma Z}=-1.78$$

(we only considered the cases for $A_{LR}^{(5)(\mathrm{SE})}(q^2)$ at a 500 GeV linear collider).

Starting from the previous expressions, Eqs. (17)–(19), it is now straightforward to calculate various kinds of contributions of self-energy type, in particular that coming from a model that implies the existence of a couple of strongly coupled vector (V) and axial vector (A) resonances. For the latter ones we shall follow the same notation as in Ref. [3], adopting the simplest treatment based on a δ-function approximation (but keeping in mind the discussion given there on the possibility of using a more realistic description without changing the essential results, i.e., the mass limits; illustrations in Figs. 1–3 are actually made with this description). We shall not abandon at this stage the customary assumption of isospin and parity conservation. Thus, the imaginary parts of the various spectral functions will simply be expressed in terms of the two quantities R_{VV}, R_{AA} with

$$R_{VV,AA}=12\pi^2F_{V,A}^2\delta(s-M_{V,A}^2)\ . \qquad (20)$$

Our investigation now proceeds in two steps. First, we assume as we did in Ref. [3] the validity of the two Weinberg sum rules (but we only fully exploited the consequences of the second one) and we made use of the experimental constraint on the parameter S, which can be written to quite a reasonable approximation as

$$-1.5\le S\le 0.5 \qquad (21)$$

only considering the *positive* upper bound. Then we combined the previous *ansätze* with the request that the experimental accuracies on $R^{(5)}$, $A_{LR}^{(5)}$, and $R_{b,\mu}$ are of a relative 1%, 1%, and 2% respectively [9] and imposed the consequent "observability" limits.

Figure 1 shows the results of our analysis for the case $\sqrt{q^2}=500$ GeV. The different curves correspond to the various observables, and the shaded area corresponds to the combined overall mass bound.

From inspection of Fig 1, the following main conclusions may be derived: (a) the only hadronic observable which contributes appreciably the bound is A_{LR}^h, which allows improving the pure leptonic result by approximately 150 GeV; (b) the resulting bounds on M_V, M_A are located in the TeV range and are rather strongly correlat-

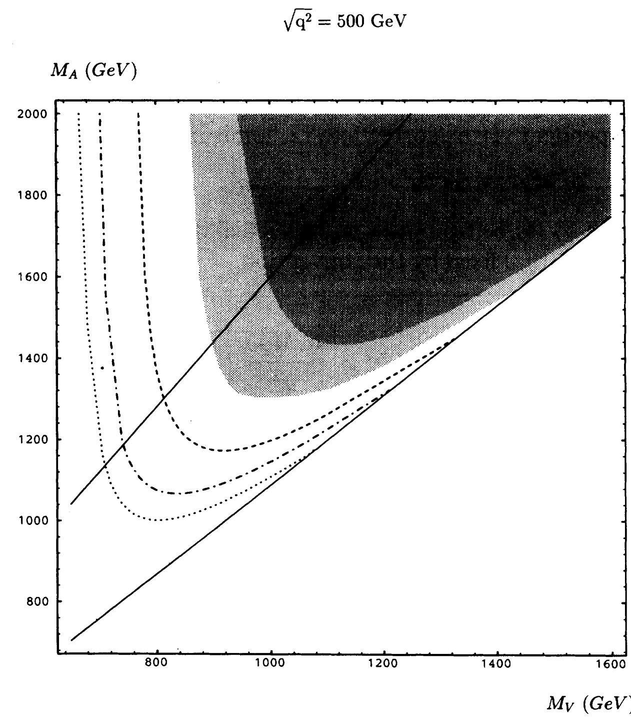

FIG. 1. Limits on M_A at variable M_V obtained at $\sqrt{q^2}=500$ GeV from σ_μ (dotted), $A_{LR,h}$ (dot-dashed) and A_τ (dashed), using the Weinberg sum rules and the experimental information on S. The lighter shaded domain represents the result of combining quadratically the two leptonic limits. The darker one corresponds to the domain allowed by the leptonic and the hadronic limits. The two full lines correspond to $M_A=1.6M_V$ and to $M_A=1.1M_V$.

ed. For the QCD-like choice $M_A/M_V=1.6$, values of M_V up to 1 TeV would be seen.

In the previous analysis, several theoretical assumptions (or prejudices?) were enforced, on which the obtained bounds certainly depend. To try to make the interconnection between the numerical output and the theoretical input more quantitatively defined might be an interesting goal. With this aim, we considered the consequences of abandoning some of the starting ingredients of our approach. Since we would personally feel uneasy in giving up the familiar isospin and parity conservation philosophy, we began by eliminating the assumptions of validity of both Weinberg sum rules and only retained a "minimal" convergence assumption $[F_{VV}(q^2)-F_{AA}(q^2)] \sim 0$, $q^2 \to \infty$, to ensure the unsubtracted form of S. This choice has two main consequences, that of introducing another degree of freedom in the analysis and that of allowing the Peskin-Takeuchi parameter S to become negative, since one has now

$$S=\left[\frac{F_V^2}{M_V^2}-\frac{F_A^2}{M_A^2}\right] \tag{22}$$

with no special indications for its sign. Thus, the experimental constraint for S, Eq. (21), will now allow both end points of the allowed interval to be saturated.

In performing our numerical analysis, we had to solve the problem of the presence of one additional degree of freedom. We decided to proceed by retaining a "prejudice relic" in which the value of the ratio F_V/M_V was bounded by the limit

$$\frac{F_V}{M_V}=2\frac{f_\rho}{m_\rho}=\frac{1}{\sqrt{2\pi}} \tag{23}$$

i.e., twice the QCD value. Higher values of the ratio would obviously increase the mass bounds accordingly, as from Eq. (22). Then, for every choice of F_V^2/M_V^2, F_A^2/M_A^2 was allowed to saturate both limits of Eqs. (21). The final results were then plotted as in the case of Fig. 1 in the (M_V,M_A) plane. In Fig. 2 we give the results of the procedure that correspond to the choice $F_V^2/M_V^2=1/2\pi$, showing that the situation has now definitely changed with respect to Fig. 1. In particular, one sees now that the effect of releasing the validity of the Weinberg sum rules is roughly that of increasing the bounds on (M_VM_A) from the 1 TeV region to the 2 TeV region for a reasonable limitation on F_V/M_V. The effect of the hadronic observables is still to increase the mass bounds by about 150 GeV.

To complete our analysis, we examined the similar situation that would occur at $\sqrt{q^2}=190$ GeV, i.e., the near future LEP 2 energy. We proceeded as before with the experimental conditions expected by previous analyses [12]. The results that we obtained are shown in Fig. 3. As one sees LEP 2 under realistic experimental conditions would be able to reveal signals of strong resonances whose masses range up to 300–350 GeV (assuming the Weinberg sum rules) or to 400–450 GeV (releasing them). These values appear relatively low in classical TC pic-

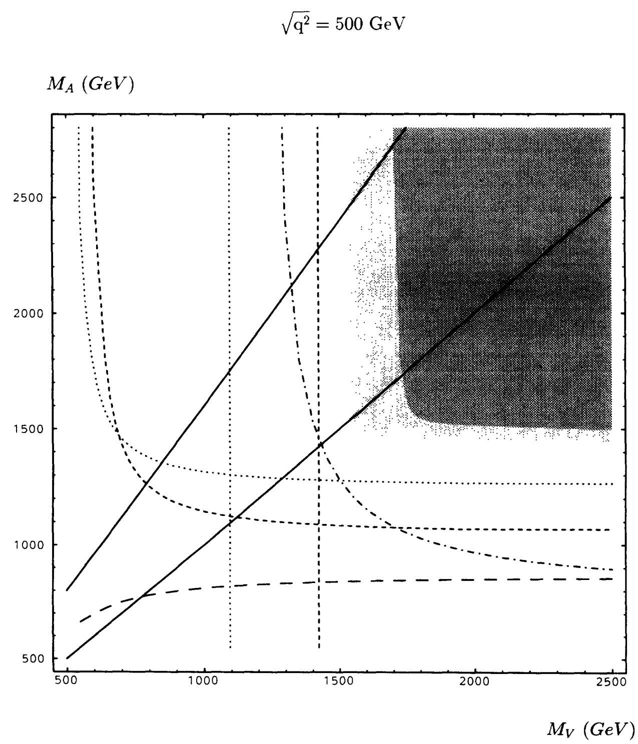

FIG. 2. Limits when releasing the Weinberg sum rules but imposing the limitation on F_V/M_V, from σ_μ (vertical, dotted), A_τ (vertical, dashed), $A_{LR,h}$ (dot-dashed), $R_{b,\mu}$ (short dashed), $R^{(5)}$ (dotted), $A_{FB,\mu}$ (long dashed). The shaded domains have the same meaning as in Fig. 1. The two full lines now correspond to $M_A=1.6M_V$ and to $M_A=M_V$.

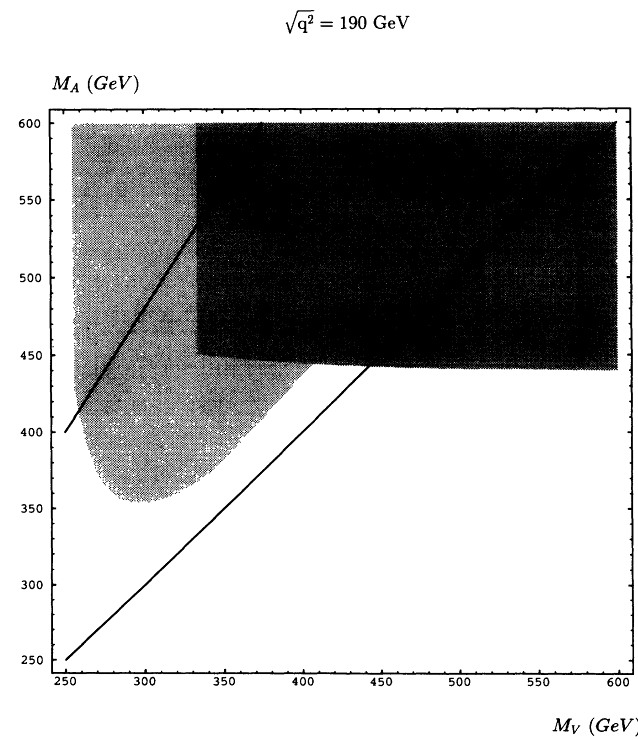

FIG. 3. Resulting domains obtained at $\sqrt{q^2}=190$ GeV (same meaning as in Fig. 2) with accuracies expected at LEP 2.

tures [13], but would certainly be much more interesting in non orthodox TC versions more recently suggested [14] implying the existence of light strongly resonant states.

In conclusion, and although our investigation was relatively qualitative, we feel that its indications should be considered as an example of the potential interest of such measurements at future e^+e^- colliders.

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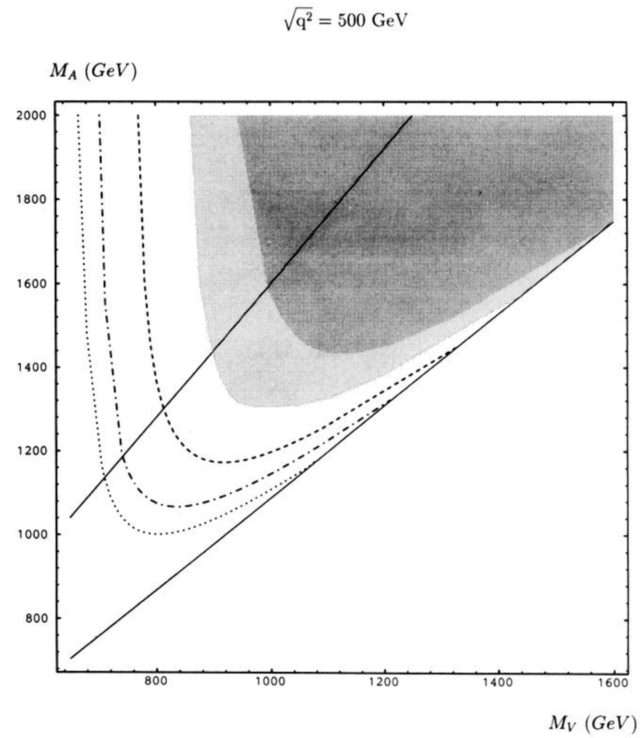

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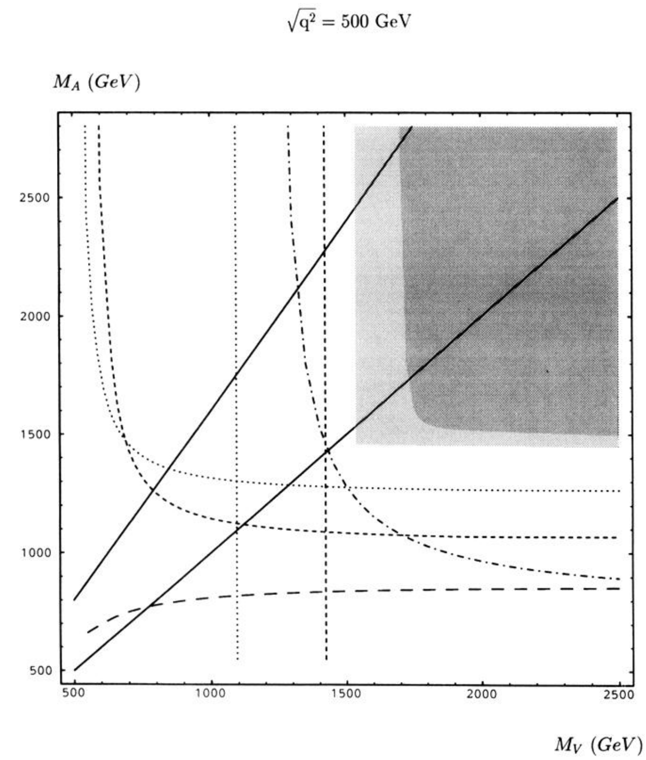

FIG. 2. Limits when releasing the Weinberg sum rules but imposing the limitation on F_V/M_V, from σ_μ (vertical, dotted), A_τ (vertical, dashed), $A_{LR,h}$ (dot-dashed), $R_{b,\mu}$ (short dashed), $R^{(5)}$ (dotted), $A_{FB,\mu}$ (long dashed). The shaded domains have the same meaning as in Fig. 1. The two full lines now correspond to $M_A = 1.6 M_V$ and to $M_A = M_V$.

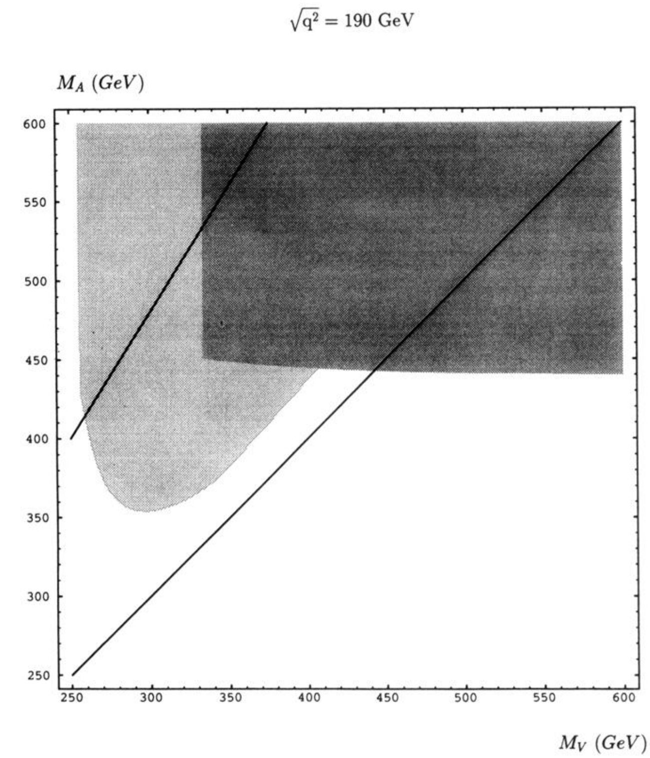

FIG. 3. Resulting domains obtained at $\sqrt{q^2}=190$ GeV (same meaning as in Fig. 2) with accuracies expected at LEP 2.