

Signals for CP Violation in Scalar Tau Pair Production at Muon Colliders

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We discuss signals for CP violation in $\mu^+\mu^- \rightarrow \tilde{\tau}_i^-\tilde{\tau}_j^+$, where $i, j = 1, 2$ label the two scalar τ mass eigenstates. We assume that these reactions can proceed through the production and decay of the heavy neutral Higgs bosons present in supersymmetric models. CP violation in the Higgs sector can be probed through a rate asymmetry even with unpolarized beams, while the CP -odd phase associated with the $\tilde{\tau}$ mass matrix can be probed only if the polarization of at least one beam can be varied. These asymmetries might be $\mathcal{O}(1)$. [S0031-9007(98)07864-8]

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In the past few years a considerable amount of effort has been devoted to investigations of the physics potential of high energy $\mu^+\mu^-$ colliders (MC) [1]. Since muons emit far less synchrotron radiation than electrons do, an MC might be significantly smaller and cheaper than an e^+e^- collider operating at the same center-of-mass energy \sqrt{s} . The main physics advantage of MC's is that the larger Yukawa coupling of muons in many cases admits copious production of Higgs bosons as s -channel resonances, allowing one to perform precision measurements of their properties [1–3]. In particular, one can search for CP violation in the couplings of Higgs bosons to heavy standard model (SM) fermions [4].

In this Letter we point out the possibility of studying CP -violating phases associated with soft supersymmetry breaking at an MC. Supersymmetry is now widely regarded to be the most plausible extension of the SM; among other things, it stabilizes the gauge hierarchy [5] and allows for the grand unification of all known gauge interactions [6]. Of course, supersymmetry must be (softly) broken to be phenomenologically viable. In general, this introduces a large number of unknown parameters, many of which can be complex. CP -violating phases associated with sfermions of the first and, to a lesser extent, second generation are severely constrained by bounds on the electric dipole moments of the electron, neutron, and muon. However, it has recently been realized [7] that cancellations between different diagrams allow some combinations of these phases to be quite large. Even in models with universal boundary conditions for soft breaking masses at some very high energy scale, the relative phase between the supersymmetric higgsino mass parameter μ and the universal trilinear soft breaking parameter A_0 can be $\mathcal{O}(1)$ [8].

Here we focus on more general models, where universality is not assumed; the phases of third generation trilinear soft breaking parameters are then essentially unconstrained. The experimental bounds on the electric dipole moments of third generation fermions are too weak to impose meaningful limits on soft breaking parameters. Phases of third generation trilinear soft breaking operators

affect the electric dipole moments of light SM fermions only at the two-loop level, e.g., through renormalization group effects [9]; since these contributions are also proportional to the respective Yukawa coupling, they will give very weak constraints on the phases of A_b and A_τ , with the possible exception of the region of very large $\tan\beta$. In fact, there is reason to believe that some of these phases might be large [10], since many proposed explanations of the baryon asymmetry of the Universe require non-SM sources of CP violation.

Unfortunately it is difficult to probe these phases through processes controlled by gauge interactions, where large CP -odd asymmetries can emerge only if some sfermion mass eigenstates are closely degenerate, with mass splitting on the order of the decay width, in which case “flavor oscillations” can occur [11–13]. On the other hand, even in the minimal supersymmetric extension of the SM, the minimal supersymmetric SM (MSSM), CP -violating phases can appear at tree level in the couplings of a single sfermion species to neutral Higgs bosons. These phases can give rise to large CP -odd asymmetries regardless of sfermion mass splittings. Here we focus on $\tilde{\tau}$ pair production. Unlike sfermions of the first two generations, $\tilde{\tau}$'s generally have sizable couplings to heavy Higgs bosons even if the latter are much heavier than M_Z . Furthermore, unlike for \tilde{b} and \tilde{t} production the charge of a produced $\tilde{\tau}$ is usually readily measurable; this is necessary for the construction of most CP -odd asymmetries. Finally, in most models sleptons are significantly lighter than squarks, making it easier to study them at lepton colliders.

Recently it has been realized [14] that CP violation in the sfermion sector will lead to mixing between the CP -even (h, H) and CP -odd (A) Higgs bosons of the MSSM. Although this is a radiative effect, it can change certain asymmetries dramatically [15]. Most of this effect is expected to come from loops involving \tilde{t} or \tilde{b} squarks. Rather than specifying the numerous free parameters of these sectors, we simply choose a value for the CP -violating $H - A$ mixing mass term $\delta m_{H,A}^2$, within the range found in Ref. [14].

For realistic $\tilde{\tau}$ masses the exchange of the lightest Higgs boson contributes negligibly to the matrix element, so that $h - A$ mixing is of little importance for us.

In general, the matrix element \mathcal{M} for $\mu^+ \mu^- \rightarrow \tilde{\tau}_i^- \tilde{\tau}_j^+$ receives contributions from γ and Z exchange as well as from the exchange of the neutral Higgs bosons of

the MSSM. The square of this matrix element for general (longitudinal or transverse) beam polarization can be computed either using standard trace techniques (employing general spin projection operators) or from the helicity amplitudes by a suitable rotation [16] from the helicity basis to a general spin basis. Both calculations give the same result,

$$\begin{aligned}
|\mathcal{M}|^2 = & \frac{e^4 |\vec{k}|^2 \sin^2 \theta}{2s} [(|V_{ij}|^2 + |A_{ij}|^2)(1 - P_L \bar{P}_L) + 2 \operatorname{Re}(V_{ij} A_{ij}^*) (P_L - \bar{P}_L) - (|V_{ij}|^2 - |A_{ij}|^2) P_T \bar{P}_T \cos(\alpha + \bar{\alpha})] \\
& + \frac{h_\mu^2}{4s} [(|P_{ij}|^2 + |S_{ij}|^2)(1 + P_L \bar{P}_L) + (|S_{ij}|^2 - |P_{ij}|^2) P_T \bar{P}_T \cos(\alpha - \bar{\alpha}) \\
& \quad + 2 \operatorname{Re}(P_{ij} S_{ij}^*) (P_L + \bar{P}_L) - 2 \operatorname{Im}(P_{ij} S_{ij}^*) P_T \bar{P}_T \sin(\alpha - \bar{\alpha})] \\
& + \frac{e^2 h_\mu |\vec{k}| \sin \theta}{\sqrt{2} s} [\operatorname{Re}(S_{ij}^* V_{ij}) (\bar{P}_L P_T \cos \alpha - P_L \bar{P}_T \cos \bar{\alpha}) + \operatorname{Im}(S_{ij}^* V_{ij}) (P_T \sin \alpha + \bar{P}_T \sin \bar{\alpha}) \\
& \quad - \operatorname{Re}(S_{ij}^* A_{ij}) (P_T \cos \alpha + \bar{P}_T \cos \bar{\alpha}) + \operatorname{Im}(S_{ij}^* A_{ij}) (P_L \bar{P}_T \sin \bar{\alpha} - \bar{P}_L P_T \sin \alpha) \\
& \quad + \operatorname{Re}(P_{ij}^* V_{ij}) (P_T \cos \alpha - \bar{P}_T \cos \bar{\alpha}) + \operatorname{Im}(P_{ij}^* V_{ij}) (P_T \bar{P}_L \sin \alpha + \bar{P}_T P_L \sin \bar{\alpha}) \\
& \quad - \operatorname{Re}(P_{ij}^* A_{ij}) (P_L \bar{P}_T \cos \bar{\alpha} + \bar{P}_L P_T \cos \alpha) + \operatorname{Im}(P_{ij}^* A_{ij}) (\bar{P}_T \sin \bar{\alpha} - P_T \sin \alpha)]. \tag{1}
\end{aligned}$$

Here, \vec{k} is the $\tilde{\tau}^-$ 3-momentum in the center-of-mass frame, θ is the scattering angle, e is the QED gauge coupling, and $h_\mu = gm_\mu / (2M_W \cos \beta)$ determines the strength of the μ Yukawa couplings, with $\tan \beta$ being the usual ratio of Higgs vacuum expectation values. V_{ij} , A_{ij} , S_{ij} , and P_{ij} are combinations of coupling factors and propagators, corresponding to vector, axial vector, scalar, and pseudoscalar couplings to $\mu^+ \mu^-$, respectively. V_{ij} and A_{ij} are dimensionless and describe γ and Z exchange (only the latter contributes to A_{ij}), while S_{ij} and P_{ij} describe Higgs exchange contributions and have a dimension of mass. Explicit expressions for these quantities will be given elsewhere [17]. Finally, P_L and \bar{P}_L are the longitudinal polarizations of the μ^- and μ^+ beams, while P_T and \bar{P}_T are the degrees of transverse beam polarization, with α and $\bar{\alpha}$ being the azimuthal angles between these polarization vectors and \vec{k} . Note that $P_L^2 + P_T^2 \leq 1$ and $\bar{P}_L^2 + \bar{P}_T^2 \leq 1$.

In this notation a CP transformation corresponds to the simultaneous exchanges $P_L \leftrightarrow -\bar{P}_L$, $P_T \leftrightarrow \bar{P}_T$, and $\alpha \leftrightarrow \bar{\alpha}$. Out of the 15 terms appearing in Eq. (1), the first five, as well as terms 8 through 11, are CP even, while the remaining six terms are CP odd. Let $C_n(i, j)$ be the coefficients of these 15 terms (bilinears in V_{ij} , A_{ij} , S_{ij} , and P_{ij} and their complex conjugates). For the coefficients multiplying CP -even factors (the first group), only the antisymmetric combinations $[C_n] \equiv C_n(1, 2) - C_n(2, 1)$ lead to CP violation through rate asymmetries. In contrast, all symmetric combinations $\{C_n\} \equiv [C_n(i, j) + C_n(j, i)]/2$ of the coefficients of the second group of terms contribute to CP -odd polarization or azimuthal angle asymmetries; these can be probed for

three different CP -even final states ($\tilde{\tau}_i^- \tilde{\tau}_i^+$, $i = 1, 2$ and the sum of $\tilde{\tau}_1^- \tilde{\tau}_2^+$ and $\tilde{\tau}_1^+ \tilde{\tau}_2^-$ production).

We emphasize that CP -odd combinations of all 15 coefficients appearing in Eq. (1) can be extracted independently, if the polarization of both beams can be controlled completely. To mention only two examples, $\{C_6\}$ can be extracted by measuring the difference of cross sections for $P_L = \bar{P}_L = +1$ and $P_L = \bar{P}_L = -1$; recall that $|P_L| = 1$ implies $P_T = 0$. $\{C_{14}\}$ can be determined by measuring $\int d\alpha d\bar{\alpha} |\mathcal{M}|^2 (\cos \alpha + \cos \bar{\alpha})$ for $P_L = \bar{P}_L = 1$, adding the same quantity for $P_T = -\bar{P}_L = 1$ (for a CP -even polarization state) and symmetrizing in the $\tilde{\tau}$ indices. In this fashion one can define nine rate asymmetries A_R and six polarization/angle asymmetries A_P ; recall that the latter can be studied for three different final states, leading to a total of 27 different asymmetries.

How many of these asymmetries can actually be measured in practice depends on the beam energy (which determines how many different final states $\tilde{\tau}_i^- \tilde{\tau}_j^+$ are accessible) and, crucially, on the extent to which the beam polarization can be controlled. If this is not possible at all, only the total rate asymmetry $\propto [C_1 + C_4]$ can be measured. If the longitudinal polarization can be tuned but $P_T = \bar{P}_T = 0$, one can in addition determine a rate asymmetry $\propto [C_2]$ (which, however, is expected to vanish, since C_2 involves only gauge interactions) and a polarization asymmetry $\propto \{C_6\}$. All other asymmetries are accessible only if at least one beam is transversely polarized. Note that asymmetries that require only one transversely polarized beam can be measured only if the azimuthal angle of the $\tilde{\tau}$'s can be reconstructed; this should be possible fairly efficiently at

least on a statistical basis, unless one is very close to the threshold (in which case the cross section is quite small anyway). Asymmetries that are accessible only if P_T and \bar{P}_T are both nonzero depend only on the difference $\alpha - \bar{\alpha}$, which is independent of \vec{k} ; this includes the polarization/angle asymmetry $\propto \{C_7\}$, which is analogous to the “production asymmetry” introduced in Ref. [4].

Some amount of longitudinal polarization will likely be present automatically, if the muons are produced from the weak decay of light mesons. This by itself is not sufficient to measure polarization asymmetries; one has to be able to tune the beam polarization, which might entail a significant reduction of the luminosity [3]. Producing transversely polarized beams will not be easy. Conventional spin rotators used for electron beams will not be effective, since the magnetic dipole moments of leptons scale as the inverse of their mass. It might nevertheless be useful to investigate what additional information might become accessible with transversely polarized beams.

To that end we present numerical results for a “typical” set of MSSM parameters: $m_A = |\mu| = |A_\tau| = 500$ GeV, gaugino mass $M_2 = 300$ GeV, $m_{\tilde{\tau}_L} = 230$ GeV, $m_{\tilde{\tau}_R} = 180$ GeV, and $\tan\beta = 10$. We set all phases to zero, except for that of A_τ which we take to be 1. The choice of M_2 affects our results only through the Higgs decay widths, which can get significant contributions from decays into neutralinos and charginos. The ratio of heavy Higgs boson masses, controlled by m_A , and soft breaking $\tilde{\tau}$ masses has been chosen such that all combinations $\tilde{\tau}_i^- \tilde{\tau}_j^+$ can be produced in the decay of on-shell Higgs bosons.

In Fig. 1 we show results for the total cross sections for $\tilde{\tau}$ pair production. In this figure we have set $H - A$ mixing to zero, but introducing a nonzero $\delta m_{H,A}^2$ in the range found in Ref. [14] has little influence on these results. The nontrivial phase between μ and A_τ leads to CP violation in the Higgs- $\tilde{\tau}$ - $\tilde{\tau}$ couplings, so that the exchange of both heavy Higgs bosons contributes to all three channels. (If CP is conserved, A exchange contributes only to $\tilde{\tau}_1^- \tilde{\tau}_2^+$ production.) However, the Higgs decay widths (~ 1.2 GeV for both A and H) are significantly larger than the $H - A$ mass difference of 400 MeV, so that only a single resonance structure is visible in the line shapes. Higgs exchange contributions completely dominate $\tilde{\tau}_1^- \tilde{\tau}_2^+$ production for $|\sqrt{s} - m_A| \leq 3$ GeV, while they are at best comparable to the gauge contributions for $\tilde{\tau}_1^- \tilde{\tau}_1^+$ production. Recall that these results are for the moderately large value $\tan\beta = 10$. Increasing $\tan\beta$ even further has little effect on the cross sections near $\sqrt{s} = m_A$, since the couplings of the heavy Higgs bosons to $\tilde{\tau}$ and μ pairs grow $\propto \tan\beta$, while the total Higgs decay widths are $\propto \tan^2\beta$, but broadens the poles by a factor $\propto \tan^2\beta$. Moreover, the Higgs exchange contributions to the matrix element scale essentially linearly in $|A_\tau|$ as long as $m_A^2 \gg M_Z^2$ and $\tan^2\beta \gg |A_\tau/\mu|$.

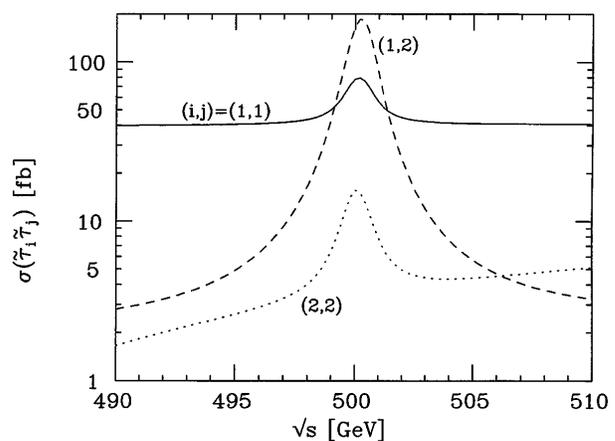


FIG. 1. Total cross sections for $\mu^+ \mu^- \rightarrow \tilde{\tau}_i^- \tilde{\tau}_j^+$. The curve labeled “(1,2)” refers to the sum of $\tilde{\tau}_1^- \tilde{\tau}_2^+$ and $\tilde{\tau}_1^+ \tilde{\tau}_2^-$ production. See the main text for the choice of parameters.

In Figs. 2 and 3 we show some “effective asymmetries,” defined as products of an asymmetry and the square root of the relevant cross section; these determine the integrated luminosity times reconstruction efficiency required to detect this asymmetry. In Fig. 2 we again set $\delta m_{H,A}^2 = 0$. In this case the total rate asymmetry $A_R(1)$ is entirely due to $h - H$ interference and is hence unmeasurably small. In contrast, near the Higgs peak the effective polarization asymmetries $\hat{A}_P(1) \propto \{C_6\}$ and $\hat{A}_P(2) \propto \{C_{12}\}$ are both very large. Recall that the former can be measured with longitudinally polarized beams, while the latter can be studied only if at least one beam is transversely polarized. In Figs. 2 and 3 we have assumed 100% beam polarization. Imperfect polarization would dilute these effective asymmetries linearly. The effective rate asymmetries $\hat{A}_R(5) \propto [C_{10}]$ and $\hat{A}_R(9) \propto [C_8]$ can also reach the level of $1 \text{ fb}^{1/2}$. Note that the latter goes through zero at $\sqrt{s} = m_H$ and falls only slowly away from the pole

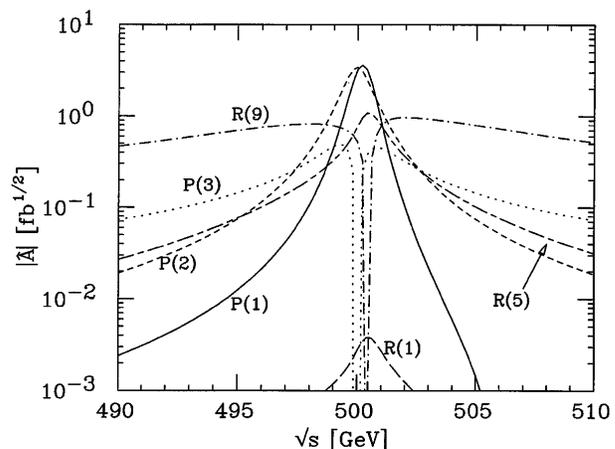


FIG. 2. Absolute values of selected asymmetries times square root of the cross section for the same set of parameters as in Fig. 1. The labels “R” and “P” refer to rate and polarization/azimuthal angle asymmetries, respectively.

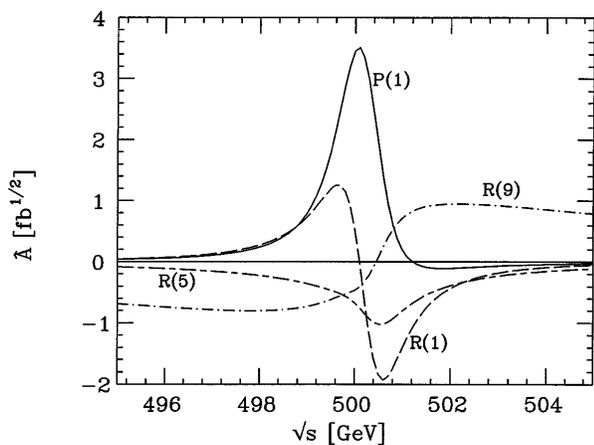


FIG. 3. Selected effective asymmetries for the same parameters as in Fig. 1, except that we have introduced a nonzero $\delta m_{H,A}^2 = 100 \text{ GeV}^2$.

region. However, this asymmetry is measurable only with one longitudinally and one transversely polarized beam; the effective asymmetry therefore scales like the square of the overall degree of beam polarization, which means that the luminosity required to see an effect scales as the inverse fourth power of the degree of polarization. Finally, the effective polarization asymmetry $\hat{A}_P(3) \propto \{C_{15}\}$ goes through zero at $\sqrt{s} = m_A$. It drops off less quickly away from the pole region than the other polarization asymmetries do, but for this set of parameters it always stays below $0.5 \text{ fb}^{1/2}$. All polarization/angle asymmetries in Figs. 2 and 3 refer to $\tilde{\tau}_1^+ \tilde{\tau}_1^-$ production; in some cases the corresponding asymmetries for $\tilde{\tau}_1 \tilde{\tau}_2$ production are even larger.

In Fig. 3 we show results for $\delta m_{H,A}^2 = 100 \text{ GeV}^2$ [14]. This increases the mass splitting between the two heavy Higgs bosons by less than 50 MeV. However, the contributions of these two bosons can now interfere, since each of them has both scalar and pseudoscalar couplings to muons. This can lead to a large effective rate asymmetry $\hat{A}_R(1)$. Note that this asymmetry goes through zero at a value of \sqrt{s} between the masses of the two Higgs eigenstates. In contrast, the effective rate asymmetry $\hat{A}_R(9)$ remains almost the same as before; in particular, it still goes through zero for \sqrt{s} very close to the mass of the heavier (mostly CP -even) Higgs boson. The measurement of these asymmetries as a function of the beam energy could therefore allow one to determine the mass splitting between the two heavy Higgs bosons; for the chosen example, studies of the overall line shapes would probably only allow to give an upper limit on this quantity.

For $\sqrt{s} \approx m_A$ the effective polarization asymmetry $\hat{A}_P(1)$ remains as in Fig. 2. It does change away from the poles; in particular, its zero moves from ~ 506

$\sim 501.2 \text{ GeV}$. However, the location of this zero will be difficult to determine, since $|\hat{A}_P(1)|$ remains small for larger \sqrt{s} . Finally, $\hat{A}_P(2)$ and $\hat{A}_P(3)$ (not shown) remain essentially the same as for $\delta m_{H,A}^2 = 0$.

In the simple example used here, measurements with unpolarized beams and with (at least) one longitudinally polarized beam are sufficient to determine the values of the two fundamental CP -violating parameters, $\delta m_{H,A}^2$ and the relative phase between μ and A_τ . Recall, however, that we have assumed that even $\tilde{\tau}_1 \tilde{\tau}_2$ can be produced in on-shell Higgs boson decays; if this is not true, all rate asymmetries will be very small. Also, $\delta m_{H,A}^2$ could itself be complex [14], if \tilde{t} or \tilde{b} pairs can be produced in Higgs decays; this would increase the number of CP -odd parameters by one. Finally, many additional CP phases can be introduced if there is nontrivial flavor mixing among sleptons [12]. In all these cases the availability of transversely polarized beams is essential to fully disentangle the sources of CP violation. We will discuss these issues in more detail in a future publication.

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