CP violation in two-meson Tau decays induced by heavy new physic

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ACP tension

The theoretical value and experimental measurement of the CP asymmetry in the decay widht of $\tau \to K_s \pi \nu_\tau$ are given respectively by [2, 6]

$$A_{CP}^{th} = 3.32(6) \times 10^{-3} \quad (1)$$

$$A_{CP}^{exp} = -3.6(2.3)(1.1) \times 10^{-3} \quad (2)$$

this sign difference implies a 2.8σ discrepancy.

An early approach to solve the tension

To get CPV, interference between amplitudes is needed

$$\mathcal{A}_j = |\mathcal{A}_j| \, e^{i\delta_j^s} e^{i\delta_j^w} \tag{3}$$

$$A_{CP} \propto |\mathcal{A}_1 + \mathcal{A}_2| - |\bar{\mathcal{A}}_1 + \bar{\mathcal{A}}_2| = -4\, |\mathcal{A}_1|\, |\mathcal{A}_2| \sin\delta^s \sin\delta^w, \tag{4}$$

with relative weak and strong phases given by

$$\delta^{s(w)} = \delta_1^{s(w)} - \delta_2^{s(w)}. \tag{5}$$

Relying on this simple fact, Ref. [3], considered heavy new physics effects in the form of interactions between tensor (NP) and SM currents to explain the anomaly. Later on, effective field theory methods challenged this possibility.

Low-energy Effective Field Theory Computation

The most general Fermi-like effective Lagrangian for the $\tau^- \to \overline{u} D \nu_\tau (D = d, s)$ decays is given by [4]

$$\mathcal{L}_{EFT} = -\frac{G_F^0 V_{uD}}{\sqrt{2}}(1 + \epsilon_L^D + \epsilon_R^D)\{\bar{\tau}\gamma_\mu(1-\gamma_5)\nu_\tau \cdot \bar{u}[\gamma^\mu - (1 - 2\hat{\epsilon}_R^D)\gamma^\mu\gamma_5] D + \bar{\tau}(1-\gamma_5)\nu_\tau \cdot \bar{u}[\hat{\epsilon}_S^D - \hat{\epsilon}_P^D\gamma_5]D + 2\hat{\epsilon}_T^D\bar{\tau}\sigma_{\mu\nu}(1-\gamma_5)\nu_\tau \cdot \bar{u}\sigma^{\mu\nu}D\} + \mathrm{h.c.} \quad (6)$$

with corresponding decay amplitude for the $\tau^- \to K^-\pi^0\nu_\tau$ decays (equations of motion were applied)

$$\mathcal{M}(\tau^- \to K^-\pi^0\nu_\tau) = \frac{G_F V_{us}}{2}\left[L_\mu H^\mu + \hat{\epsilon}_S^* L H + 2\hat{\epsilon}_T^* L_{\mu\nu} H^{\mu\nu}\right], \quad (7)$$

where the lepton and hadron currents are defined as

$$\begin{aligned} L &= \bar{u}(p_{\nu_\tau})(1+\gamma_5)u(p_\tau)\,, && (8)\\ L_\mu &= \bar{u}(p_{\nu_\tau})\gamma_\mu(1-\gamma_5)u(p_\tau)\,, && (9)\\ L_{\mu\nu} &= \bar{u}(p_{\nu_\tau})\sigma_{\mu\nu}(1+\gamma_5)u(p_\tau)\,, && (10)\end{aligned}$$

and

$$\begin{aligned} H &= \left\langle \pi^0(p_\pi)K^-(p_K)|\bar{s}d|0\right\rangle = \frac{\Delta_{K\pi}}{m_s-m_u}F_0(s)\,, && (11)\\ H^\mu &= \left\langle \pi^0(p_\pi)K^-(p_K)|\bar{s}\gamma^\mu d|0\right\rangle \\ &= \left[(p_\pi-p_K)^\mu+\frac{\Delta_{K\pi}}{s}q^\mu\right]F_+(s)-\frac{\Delta_{K\pi}}{s}q^\mu F_0(s)\,, && (12)\\ H^{\mu\nu} &= \left\langle \pi^0(p_\pi)K^-(p_K)|\bar{s}\sigma^{\mu\nu} d|0\right\rangle \\ &= iF_T(s)(p_\pi^\mu p_K^\nu - p_K^\mu p_\pi^\nu)\,. && (13)\end{aligned}$$

From the latter matrix elements it is straightforward to obtain the *CP-asymmetry rate* and the *Forward-Backward asymmetry* [8]

$$A_{CP}^{\rm rate}(\tau^{\pm}\to K^{\pm}\pi^{0}\nu_{\tau}) = \frac{\Gamma(\tau^{+}\to K^{+}\pi^{0}\bar{\nu}_{\tau}) - \Gamma(\tau^{-}\to K^{-}\pi^{0}\nu_{\tau})}{\Gamma(\tau^{+}\to K^{+}\pi^{0}\bar{\nu}_{\tau}) + \Gamma(\tau^{-}\to K^{-}\pi^{0}\nu_{\tau})}$$

$$= \frac{\Im m[\hat{\epsilon}_T] G_F^2 |V_{us}|^2 S_{EW}}{128\pi^3 M_\tau^2 \Gamma(\tau\to K\pi^0\nu_\tau)} \int_{(m_K+m_\pi)^2}^{M_\tau^2} \mathrm{d}s \left(1-\frac{M_\tau^2}{s}\right)^2$$

$$\lambda^{3/2}(s, m_K^2, m_\pi^2) \times |F_T(s)||F_+(s)| \sin[\delta_T(s)-\delta_+(s)]\,. \qquad (14)$$

The constraints on $\delta_T(s) - \delta_+(s)$ (given by Watson's theorem) and on $\Im m[\hat{\epsilon}_T]$ (from $D - \bar{D}$ mixing and the neutron edm) preclude a natural explanation of the BaBar asymmetry by *heavy new physics* [4].

CP Observables, $K_0 - \bar{K}_0$ *mixing*

The expression in the second and third lines of eq. (14) corresponds to the new physics (NP) contribution to A_{CP}^{rate}, which coincides with A_{CP}^{rate} for those modes without neutral Kaons. In the $\tau^{\pm} \to K^{\pm} K_S \nu_\tau$ case, this observable is dominated by the SM contribution, according to [3]

$$A_{CP}^{rate} = \frac{A_{CP}^{rate}|_{SM} + A_{CP}^{rate}|_{NP}}{1 + A_{CP}^{rate}|_{SM} \times A_{CP}^{rate}|_{NP}} \, . \tag{15}$$

CP Observables, *FB asymmetry*

$$A_{FB}^{\tau^- \to K^- \pi^0 \nu_\tau}(s) = \frac{\int_0^1 \frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha - \int_{-1}^0 \frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha}{\int_0^1 \frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha + \int_{-1}^0 \frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha}, \quad (16)$$

The *CP-FB asymmetry* actually fulfills the relation

$$A_{FB}(s) = 3/2 < \cos\alpha > (s) \quad (17)$$

with

$$< \cos\alpha > (s) = \frac{\int_{-1}^1 \cos\alpha \left(\frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha \right)}{\int_{-1}^1 \left(\frac{\mathrm{d}^2\Gamma(\tau^- \to K^- \pi^0 \nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha \right)} = \frac{N(s)}{D(s)}. \quad (18)$$

Where, in turn

$$
\begin{aligned}
N(s) &= -\frac{4}{3}\Delta_{K\pi}\lambda^{1/2}(s, m_K^2, m_\pi^2)\Re e\left[\left(1+\frac{\hat{\epsilon}_S s}{M_\tau(m_s-m_u)}\right)F_+(s)F_0^*(s)\right] \\
&+ \frac{8s}{3M_\tau}\Delta_{K\pi}\lambda^{1/2}(s, m_K^2, m_\pi^2)\Re e\left[\hat{\epsilon}_T^*\left(1+\frac{\hat{\epsilon}_S s}{M_\tau(m_s-m_u)}\right)F_T(s)F_0^*(s)\right], \\
D(s) &= \frac{2}{3}\lambda(s, m_K^2, m_\pi^2)\left(1+\frac{2s}{M_\tau^2}\right)|F_+(s)|^2 + 2\Delta_{K\pi}^2\left|1+\frac{\hat{\epsilon}_S s}{M_\tau(m_s-m_u)}\right|^2 \\
&|F_0(s)|^2 + \frac{8}{3}\lambda(s, m_K^2, m_\pi^2)\left[s|\hat{\epsilon}_T|^2\left(2+\frac{s}{M_\tau^2}\right)|F_T(s)|^2\right. \\
&- \frac{3s}{M_\tau}\Re e[\hat{\epsilon}_T F_+(s)F_T^*(s)].
\end{aligned}
$$

(19)

CP observables, a *FOM for CP-FBa*

Since the FB asymmetry observable is not given by just a number, as the rate CP asymmetry (*CP-FBa*), then we propose the following *figure of merit* to represent the *CPV signal strength*

$$F(\hat{\epsilon}_S^D, \hat{\epsilon}_T^D) = \int_{(m_1+m_2)^2}^{M_\tau^2} \left| A_{FB}(s; \hat{\epsilon}_S^D, \hat{\epsilon}_T^D) \right| \quad ds \,. \tag{20}$$

The optimal choice maximizing this quantity, shall be considered the maximal *CP-FBa*.

Form-factors input

We will use a dispersive representation for the form factors.
For the pion vector form factor we used a thrice subtracted representation and one subtraction for the scalar and tensor ones. [15, 5, 10, 14]

$$F_+^{\pi\pi}(s) = \exp\left[\alpha_1 s + \frac{\alpha_2}{2} s^2 + \frac{s^3}{\pi}\int_{4m_\pi^2}^{\infty} \frac{\mathrm{d}s'}{s'^3} \frac{\delta_+(s')}{s' - s - i0}\right], \tag{21}$$

$$F_T^{\pi\pi}(s) = \exp\left[\frac{s^3}{\pi}\int_{4m_\pi^2}^{\infty} \frac{\mathrm{d}s'}{s'} \frac{\delta_T(s')}{s' - s - i0}\right], \tag{22}$$

$$F_0^{\pi\pi}(s) = \exp\left[\frac{s^3}{\pi}\int_{4m_\pi^2}^{\infty} \frac{\mathrm{d}s'}{s'} \frac{\delta_0(s')}{s' - s - i0}\right]. \tag{23}$$

Tensor form factor phase

As we will be seeing, Watson's final state theorem limits our possibilities to observe CPV induced by heavy new physics in these decays. Due to these restrictions we (also [4, 5, 8]) estimate the phase of the tensor form factor considering

$$\delta_+(s) - \delta_T(s) = 2\delta_+^{inel}(s) \tag{24}$$

Where $\delta_+^{inel}(s)$ is given by the inelastic effects arising from the pion vector form factor only.

Form-factors input

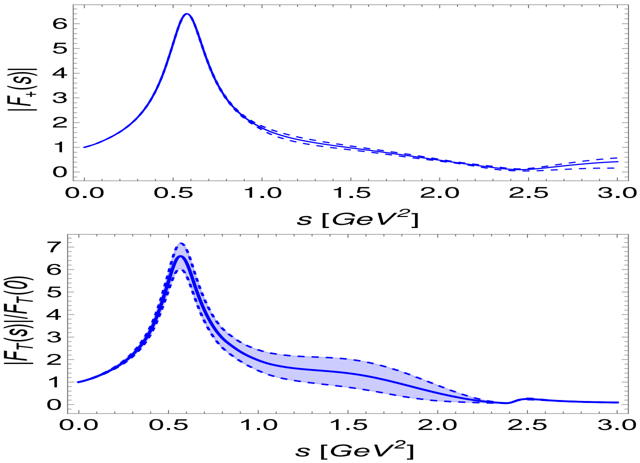

Figure: On top, pion vector form factor. At the bottom, tensor form factor derived from the estimator described in the latter section.

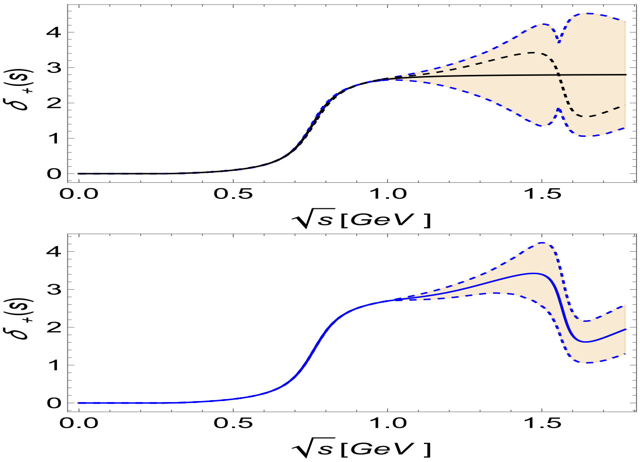

Figure: On top, tensor form factor phase. At the bottom, vector form factor phase.

Watson's final state theorem

Let us consider the matrix element $\langle(\pi\pi)_I; out|J|\Omega\rangle$, where J is some current with the property of being time reversal invariant $\mathbf{T}J\mathbf{T}^{-1} = J$, then

$$
\begin{aligned}
\langle(\pi\pi)_I; out|J|\Omega\rangle \quad &= \left(\langle(\pi\pi)_I; in|J|\Omega\rangle\right)^* \\
&= \sum_n \left(\langle(\pi\pi)_I; in|n; out\rangle\langle n; out|J|\Omega\rangle\right)^* \\
&= \langle(\pi\pi)_I; out|(\pi\pi)_I; in\rangle \left(\langle(\pi\pi)_I; out|J|\Omega\rangle\right)^* \\
&= e^{2i\delta_I} \left(\langle(\pi\pi)_I; out|J|\Omega\rangle\right)^* , \qquad (25)
\end{aligned}
$$

Thus, in the elastic zone, the phase of the strong amplitude depends only on the quantum numbers involved in the elastic process $(\pi\pi)_I \to (\pi\pi)_I$, which coincides with the one from scattering phase shift.

Bounds on the New Physics coefficients

The Fermi-like Lagrangian is obtained as the relevant part of the low-energy limit of the SMEFT Lagrangian. This connection is useful to constrain the imaginary part of the scalar and tensor Wilson coefficients We will start with the tensor operators. In the gauge basis one has (see [4, 8] for more details)

$$\begin{aligned}\mathcal{L}_{SMEFT} &\supset [C^{(3)}_{\ell equ}]_{klmn}(\bar{\ell}^i_{Lk}\sigma_{\mu\nu}e_{Rl})\epsilon^{ij}(\bar{q}^j_{Lm}\sigma^{\mu\nu}u_{Rn}) + \mathrm{h.c.}\\ &= [C^{(3)}_{\ell equ}]_{klmn}[(\bar{\nu}_{Lk}\sigma_{\mu\nu}e_{Rl})(\bar{d}_{Lm}\sigma^{\mu\nu}u_{Rn}) - (\bar{e}_{Lk}\sigma_{\mu\nu}e_{Rl})(\bar{u}_{Lm}\sigma^{\mu\nu}u_{Rn})]\\ &+\mathrm{h.c.},\end{aligned} \tag{26}$$

where the left-handed lepton and quark $SU(2)_L$ doublets are $\ell_L = (\nu_L, e_L)^T$ and $q_L = (u_L, d_L)^T$ and the corresponding singlets are e_R and u_R. $SU(2)_L$(family) indices are $i, j(k, l, m, n)$, respectively. In the *down* mass basis,

$$\begin{aligned}\mathcal{L}_{SMEFT} \quad &\supset [C^{(3)}_{\ell equ}]_{klmn}[(\bar{\nu}_{Lk}\sigma_{\mu\nu}e_{Rl})(\bar{d}_{Lm}\sigma^{\mu\nu}u_{Rn})\\ &-V_{am}(\bar{e}_{Lk}\sigma_{\mu\nu}e_{Rl})(\bar{u}_{La}\sigma^{\mu\nu}u_{Rn})] + \mathrm{h.c.}\end{aligned} \tag{27}$$

Bounds on the Wilson coefficients

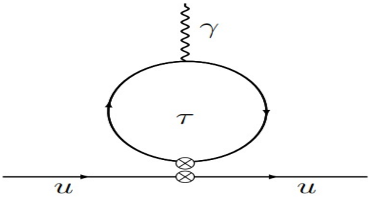

Figure: Feynman diagrams involved in the EDM and $D_0 - \bar{D}_0$ mixing.

The tensor operator $(\overline{\tau}_L\sigma_{\mu\nu}\tau_R)(\overline{u}_L\sigma^{\mu\nu}u_R)$ contributes to the *up-quark EDM* $d_u(\mu)$ through the left diagram, with the interaction term

$$\mathcal{L}_D = -\frac{i}{2}d_u(\mu)\overline{u}\sigma_{\mu\nu}\gamma_5 uF^{\mu\nu}, \qquad (28)$$

which relates to the neutron EDM via the flavour-diagonal tensor charge [22, 23, 24]

$$d_n = g_T^u(\mu)d_u(\mu). \qquad (29)$$

Bounds on the Wilson coefficients

After a double insertion of the operator $(\overline{\tau_L}\sigma_{\mu\nu}\tau_R)(\overline{c}_L\sigma^{\mu\nu}u_R)$ into the right diagram, one arrives at the effective Hamiltonian, in the $m_c = m_u = 0$ limit

$$\mathcal{H}_{eff}^{C=2} = C_2^{'}(\overline{c}_L^{\alpha}u_R^{\alpha})(\overline{c}_L^{\beta}u_R^{\beta}) + C_3^{'}(\overline{c}_L^{\alpha}u_R^{\beta})(\overline{c}_L^{\beta}u_R^{\alpha}), \tag{30}$$

where the coefficients, the off-diagonal matrix elements of the latter Hamiltonian, and the theoretical mixing parameters of the *D-meson* system given respectively by

$$C_2^{'} = \frac{1}{2}C_3^{'} = 16G_F^2\frac{m_\tau}{\pi}\left(V_{ud}V_{cd}\left[\epsilon_T\right]_{3311} + \left(V_{us}V_{cs}\left[\epsilon_T\right]_{3321}\right)^2 log\frac{\Lambda}{\mu_\tau}\right., \tag{31}$$

$$M_{12}^{NP} = \frac{1}{M_D}\left[C_2\langle D_0|(\overline{c}_L^{\alpha}u_R^{\alpha})(\overline{c}_L^{\beta}u_R^{\beta})|\bar{D}\rangle + C_3\langle D_0|(\overline{c}_L^{\alpha}u_R^{\beta})(\overline{c}_L^{\alpha}u_R^{\beta})|\bar{D}_0\rangle\right], \tag{32}$$

$$x_{12}^{NP} = \frac{2|M_{12}^{NP}|}{\Gamma_D}, \qquad \phi_{12}^{NP} = arg\left(\frac{M_{12}^{NP}}{\Gamma_{12}}\right), \tag{33}$$

Imaginary Part bounds

The LEFT (Low-Energy EFTor Fermi-like) and SMEFT coefficients are related at tree level by $[C^3_{\ell equ}]_{33mn} = -2\sqrt{2} G_F V_{uD} \hat{\epsilon}^{D*}_T$.
(see [8] for further details)

- If the neutron EDM [1] is mainly contributed by a single $\hat{\epsilon}^D_T$, then

$$d_u(\mu) = -2\sqrt{2} G_F \frac{eM_\tau}{\pi^2} V^2_{uD} \Im m[\hat{\epsilon}^D_T(\mu)] \log \frac{\Lambda}{\mu} , \quad (34)$$

which yields $|\Im m[\hat{\epsilon}^s_T]| \lesssim 4 \times 10^{-6}$ and $|\Im m[\hat{\epsilon}^d_T]| \lesssim 8 \times 10^{-5}$ for $\Lambda \gtrsim 100$ GeV and $\mu = 2$ GeV. The constraint on $D^0 - \bar{D}^0$ mixing will then come from $\phi = \Im m \left[V_{uD} V_{cD} \epsilon^D_{T(S)} \right]$.

[1] We recall that [8] $d_n = g^u_T(\mu) d_u(\mu)$, with $g^u_T(2\,\mathrm{GeV}) = -0.204(14)$ [22, 23], $|d_n| < 1.8 \times 10^{-26} e$ cm at 90% confidence level, C. L. [24].

Imaginary Part bounds

We turn now to the scalar operators in the SMEFT , which, in the fermion mass basis read

$$\mathcal{L}_{SMEFT} \supset [C^{(1)}_{\ell equ}]_{klmn} \left[(\bar{\nu}_{Lk} e_{Rl}) \left(\bar{d}_{Lm} u_{Rn} \right) - V_{am} (\bar{e}_{Lk} e_{Rl})(\bar{u}_{La} u_{Rn}) \right] + [C_{\ell edq}]_{klmn} \left[V^*_{an} (\bar{\nu}_{Lk} e_{Rl}) \left(\bar{d}_{Rm} u_{La} \right) + (\bar{e}_{Lk} e_{Rl}) \left(\bar{d}_{Rm} d_{Ln} \right) \right] . \quad (35$$

LEFT and SMEFT coefficients are related by (we neglect the contribution proportional to V^*_{ub})

$$-2\sqrt{2} G_F V^*_{uD} \hat{\epsilon}^{D*}_S = [C^{(1)}_{\ell equ}]_{33m1} + V^*_{ud} [C_{\ell edq}]_{33m1} + V^*_{us} [C_{\ell edq}]_{33m2}. \quad (36)$$

In the scalar case the strongest constraints come from $D^0 - \bar{D}^0$ mixing, implying $\Im m \left[\hat{\epsilon}^D_S \right] \in [-3.1, 1.6] \times 10^{-4}$ at 95% C.L as can be seen in [8].

Real Part bounds

The real parts of the Wilson coefficients are most effectively constrained by analyzing CP-conserving inclusive and exclusive semi-leptonic tau decays. In the following we will consider only the results from the exclusive study [25]

$$\Re e[\epsilon_S^d] = \left(1.0^{+0.6}_{-3.4}\right) \times 10^{-1}\,, \quad \Re e[\epsilon_T^d] = \left(0.4^{+4.3}_{-4.6}\right) \times 10^{-2}\,, \tag{37}$$

with a correlation coefficient of 0.452, and

$$\Re e[\epsilon_S^s] = (0.8 \pm 0.9) \times 10^{-2}\,, \quad \Re e[\epsilon_T^s] = (0.5 \pm 0.8) \times 10^{-2}\,, \tag{38}$$

with a negligible correlation coefficient of -0.057.

CP asymmetries in the $\pi^{\pm}\pi^0$ channel

In this case the SM contribution is negligible, so using, for the NP part, we find

$$A_{CP}^{rate}|_{\pi\pi} \leq 3 \times 10^{-5}\,, \tag{39}$$

Clearly too tiny to be observed soon

The maximal $A_{FB}|_{\pi\pi}$ is shown by a blue solid line in the figure.

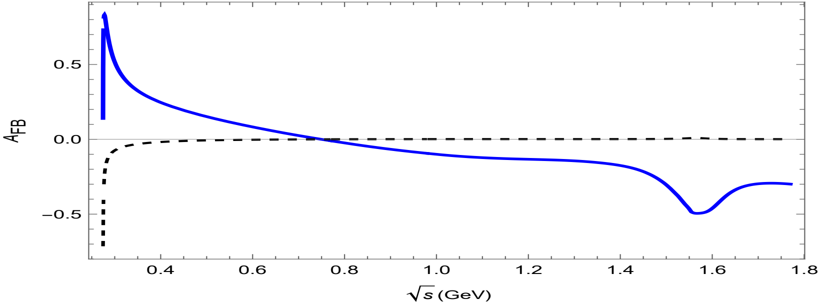

Figure: Maximal $A_{FB}^{\tau\to\pi\pi^0\nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^d] = -0.46$, $\Im m[\epsilon_S^d] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^d] = -7.2 \times 10^{-2}$ and $\Im m[\epsilon_T^d] = 8 \times 10^{-5}$, compared to the SM case (black dashed line).

FOM for the $\tau \to \pi\pi^0\nu_\tau$ channel

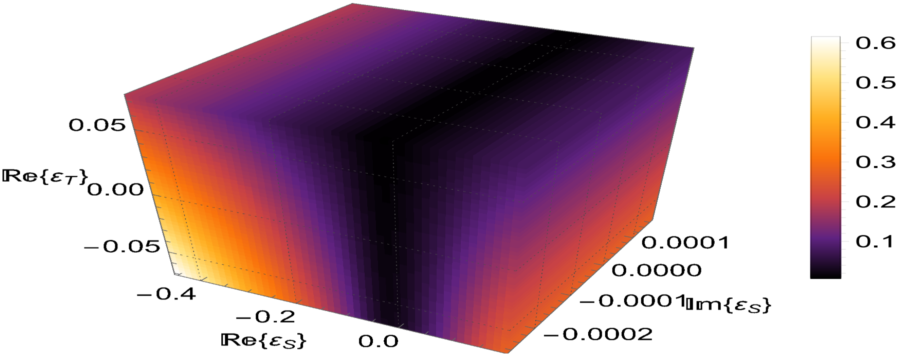

Figure: 3D density plot for the FoM on the available parameter space, setting $|\Im m[\hat{\epsilon}_T^d]| = 8 \times 10^{-5}$, for the $\pi\pi^0$ decay channel.

CP asymmetries in the $K^{\pm}K_S$ channel

The most interesting channel is given by the one with $K^{\pm}K_S$ mesons in the final state, whose total asymmetry rate (including $K_0 - \bar{K}_0$ oscillation effects in eq. 15) is given by

$$A_{CP}^{rate}|_{KK} = 3.8 \times 10^{-3}. \tag{40}$$

with a maximal NP contribution given by

$$A_{CP}^{rate}|_{KK,NP} \leq 2.3 \times 10^{-4} \tag{41}$$

Noticeably, a 5% precision on the measurement of A_{CP}^{rate} would already be sensitive to the maximum allowed NP contribution in this case.

CP asymmetries in the $K^{\pm}K_S$ channel

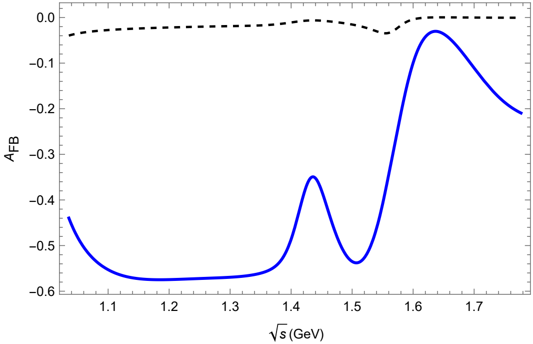

Figure: Maximal $A_{FB}^{\tau\to K_S K \nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^d] = -0.46$, $\Im m[\epsilon_S^d] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^d] = -7.2 \times 10^{-2}$ and $\Im m[\epsilon_T^d] = 8 \times 10^{-5}$, compared to the SM case (black dashed line).

CP violating asymmetries in the $K^{\pm}K_S$ channel

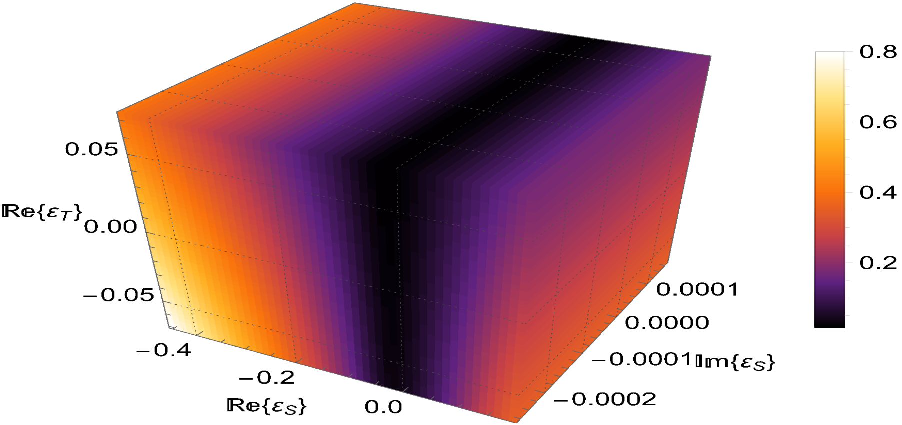

Figure: 3D density plot for the FOM described in eq. (20) on the parameter space, fixing $|\Im m[\hat{\epsilon}_T^d]| = 8 \times 10^{-5}$, for the KK_S decay channel.

CP asymmetries in the $K^{\pm}\pi^0$ channel

Finally, the maximal rate for the $K^{\pm}\pi^0$ channel is given by

$$A_{CP}^{rate}|_{K\pi} \leq 6 \times 10^{-7}\,, \tag{42}$$

Again, a too small effect to be observed soon.
The corresponding A_{FB} is displayed in the figure

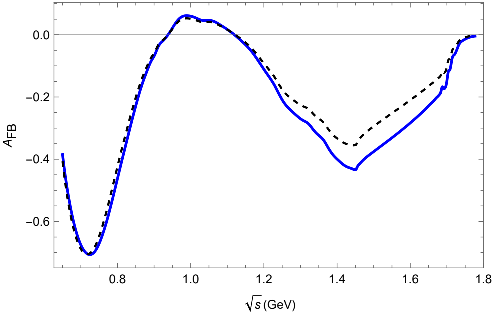

Figure: Maximal $A_{FB}^{\tau\to K\pi^0\nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^s] = 2.3 \times 10^{-2}$, $\Im m[\epsilon_S^s] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^s] = 1.8 \times 10^{-2}$ and $\Im m[\epsilon_T^s] = 4 \times 10^{-6}$, compared to the SM case (black dashed line).

FOM for the $\tau \to K\pi^0\nu_\tau$ channel

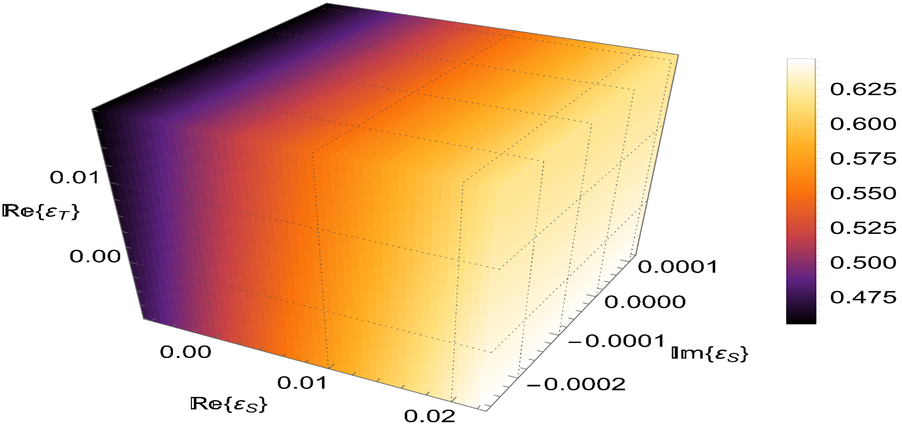

Figure: 3D density plot for the FoM described in eq. (20) on the parameter space, with $|\Im m[\hat{\epsilon}_T^d]| = 4 \times 10^{-6}$, for the decay channel with K^- and π^0 in the final state.

Conclusions

Solving the puzzle related to the anomalous BaBar measurement of A_{CP}^{rate} for the $K_S\pi$ channels, whose odd result implies a deviation at the level of a 2.8 σ, can shed light on novel CP violation mechanisms. EFTs show that heavy new physics cannot explain it unless unnatural fine-tuning is invoked.

In this work we have examined (within the same framework provided by EFT) the rate and forward-backward asymmetries for the other two-meson tau decay modes, focusing on the $(\pi/K)\pi^0$ and KK_S cases, being the most promising one turns out to be the di-Kaon channel, where measurements of A_{CP}^{rate} with a 5% precision would already be sensitive to NP contributions that are currently allowed.

We hope that our analysis motivates the Tau physics group at Belle-II to tackle these interesting analyses that we are proposing. They will also be attractive for the planned super-tau-charm factories. If future facilities produce polarized taus, this will open a bunch of new CP violating measurements exploiting this feature.

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