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

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A_{CP} tension

The theoretical value and experimental measurement of the CP asymmetry in the decay width of $\tau \to K_s \pi \nu_\tau$ are given respectively by [2, 3]

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

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

this sign difference implies a 2.8σ discrepancy. A_{CP}^{exp} is not in tension with the binned CP asymmetry (compatible with zero) measured by Belle [4]. An unexpectedly large CPV in the lepton sector can allow to understand baryogenesis via leptogenesis [5]. In case of eventual confirmation of the BaBar result, related anomalies should be seen in the $K_S \pi^{\pm} \pi^0$ and $K_S K^{\pm}$ modes.

An early approach to solve the tension

To get CPV, interference between amplitudes is needed

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

$$\begin{aligned} A_{CP} &\propto |\mathcal{A}_1 + \mathcal{A}_2|^2 - \left|\bar{\mathcal{A}}_1 + \bar{\mathcal{A}}_2\right|^2 \\ &= -4\,|\mathcal{A}_1|\,|\mathcal{A}_2| \sin\delta^s \sin\delta^w, \end{aligned} \tag{4}$$

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. [6], 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 [7, 8, 9] challenged this possibility.

Low-energy Effective Field Theory Computation

The most general Fermi-like effective Lagrangian for the $\tau^- \to \bar{u} D \nu_\tau (D = d, s)$ decays is given by [10, 7]

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

with corresponding decay amplitude for the $\tau^- \to K^-\pi^0\nu_\tau$ decays (equations of motion were applied) [1]

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

Low-energy Effective Field Theory Computation

where the lepton and hadron currents are defined as

$$L = \bar{u}(p_{\nu_\tau})(1+\gamma_5)u(p_\tau)\,, \tag{8}$$

$$L_\mu = \bar{u}(p_{\nu_\tau})\gamma_\mu(1-\gamma_5)u(p_\tau)\,, \tag{9}$$

$$L_{\mu\nu} = \bar{u}(p_{\nu_\tau})\sigma_{\mu\nu}(1+\gamma_5)u(p_\tau)\,, \tag{10}$$

and $(q^\mu = (p_\pi + p_K)^\mu,\ s = q^2,\ \Delta_{K\pi} = m_K^2 - m_\pi^2)$

$$H = \left\langle \pi^0(p_\pi)K^-(p_K)|\bar{s}d|0\right\rangle = \frac{\Delta_{K\pi}}{m_s - m_u}F_0(s)\,, \tag{11}$$

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

$$\begin{aligned} H^{\mu\nu} &= \left\langle \pi^0(p_\pi)K^-(p_K)|\bar{s}\sigma^{\mu\nu} d|0\right\rangle \\ &= iF_T(s)(p_\pi^\mu p_K^\nu - p_K^\mu p_\pi^\nu)\,. \end{aligned} \tag{13}$$

CP Observables, *CP asymmetry rate*

From the latter matrix elements it is straightforward to obtain the *CP-asymmetry rate* and the *Forward-Backward asymmetry* [11, 12]

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

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

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

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* [7].

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 [6]

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

CP Observables, *CP-FB asymmetry*

Another observable we study[1] is given by the *Forward-Backward-CP* asymmetry, which comprises the effects coming from the time-dependent mixing of the K_S and its subsequent *CP-forbidden* decay to a dipion (α is the angle between τ and a meson -typically the charged one- in the hadron rest frame [13])

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

$$\frac{\mathrm{d}^2\Gamma(\tau^-\to K_S[\pi\pi]\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha} = \frac{\mathrm{d}^2\Gamma(\tau^-\to \overline{K}^0\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\Gamma\left(\overline{K}^0(t)\to\pi\pi\right), \tag{16}$$

$$\frac{\mathrm{d}^2\Gamma(\tau^+\to K_S[\pi\pi]\pi^+\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha} = \frac{\mathrm{d}^2\Gamma(\tau^+\to K^0\pi^+\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\Gamma\left(K^0(t)\to\pi\pi\right).$$

CP Observables, *FB asymmetry*

$$A_{FB}^{\tau^-\to K_S\pi^-\nu_\tau}(s) = \frac{\int_0^1 \frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha - \int_{-1}^0 \frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha}{\int_0^1 \frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha + \int_{-1}^0 \frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha}, \tag{17}$$

The *CP-FB asymmetry* actually fulfills the relation

$$A_{FB}(s) = 3/2 < \cos\alpha > (s) \tag{18}$$

with

$$< \cos\alpha > (s) = \frac{\int_{-1}^1 \cos\alpha \left(\frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha\right)}{\int_{-1}^1 \left(\frac{\mathrm{d}^2\Gamma(\tau^-\to K_S\pi^-\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha\right)} = \frac{N(s)}{D(s)}. \tag{19}$$

CP Observables, *FB asymmetry*

The previous nominator and denominator are

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

CP observables

To calculate $A_{CP-FB}^{\tau\to K_S\pi\nu_\tau}(s)$, we use [1] the approximate relation

$$A_{CP-FB}^{\tau\to K_S\pi\nu_\tau}(s) \approx \quad (<\cos\alpha>^{\tau_-} + <\cos\alpha>^{\tau_+})A_{CP}^K + <\cos\alpha>^{\tau_-} - <\cos\alpha>^{\tau_+}$$

where

$$<\cos\alpha>^{\tau_-} \pm <\cos\alpha>^{\tau_+} = \int_0^1 \frac{\mathrm{d}^2\Gamma(\tau^-\to K^0\pi\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha \pm \int_{-1}^0 \frac{\mathrm{d}^2\Gamma(\tau^+\to \overline{K}^0\pi\nu_\tau)}{\mathrm{d}s\,\mathrm{d}\cos\alpha}\mathrm{d}\cos\alpha$$

and $A_{CP}^K = A_{CP}^{K\to\pi\pi}(t_1 << \Gamma_S^{-1}, \Gamma_S^{-1} << t_2 << \Gamma_L^{-1}) \sim -3.32\cdot 10^{-3}$ is nothing but the time-dependent difference between the production rate of the *di-pion* from the neutral Kaon.

CP observables, a *FOM for CP-FB*

Since the FB asymmetry observable is not given by just a number, as the rate CP asymmetry (*CP-FB*), then we propose [1] the following *figure of merit* to represent the *CPV signal strength*

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

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

Form-factors input

We use [1] 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 (similarly for the other channels, known with less precision) [14, 15, 8, 16, 17, 18, 19, 20, 21].

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

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

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

Tensor form factor phase

As we showed [1], Watson's final state theorem limits our possibilities to observe CPV induced by heavy new physics in these decays. Due to these restrictions we [1] (see also [7, 8, 12]) estimate the phase of the tensor form factor considering

$$\left|\delta_+(s) - \delta_T(s)\right| = 2\delta_+^{\text{inel}}(s)\,, \tag{24}$$

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

Bounds on the Wilson coefficients

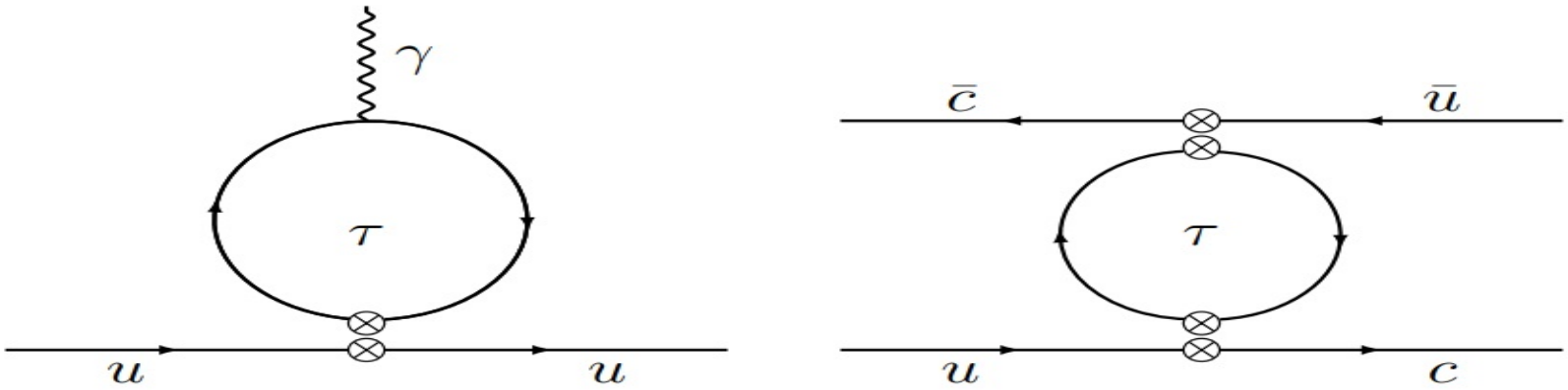

Figure: Feynman diagrams involved in the EDM and $D^0 - \bar{D}^0$ mixing.

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

$$\mathcal{L}_D = -\frac{i}{2} d_u(\mu) \overline{u} \sigma_{\mu\nu} \gamma_5 u F^{\mu\nu}, \tag{25}$$

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

$$d_n = g_T^u(\mu) d_u(\mu). \tag{26}$$

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

$$\mathcal{H}_{eff}^{C=2} = C_2^{'}(\overline{c}_L^{\alpha}u_R^{\alpha})(\overline{c}_L^{\beta}u_R^{\beta}) + C_3^{'}(\overline{c}_L^{\alpha}u_R^{\beta})(\overline{c}_L^{\beta}u_R^{\alpha}), \tag{27}$$

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

$$C_2^{'} = \frac{1}{2}C_3^{'} = 16G_F^2\frac{m_\tau^2}{\pi^2}\left(V_{ud}V_{cd}\left[\epsilon_T\right]_{3311} + V_{us}V_{cs}\left[\epsilon_T\right]_{3321}\right)^2 log\frac{\Lambda}{\mu_\tau}, \tag{28}$$

$$M_{12}^{NP} = \frac{1}{2M_D}\left[C_2'(\mu)\langle D^0|(\overline{c}_L^{\alpha}u_R^{\alpha})(\overline{c}_L^{\beta}u_R^{\beta})|\bar{D}^0\rangle(\mu) + C_3'(\mu)\langle D^0|(\overline{c}_L^{\alpha}u_R^{\beta})(\overline{c}_L^{\beta}u_R^{\alpha})|\bar{D}^0\rangle(\mu)\right] \tag{29}$$

$$x_{12}^{NP} = \frac{2|M_{12}^{NP}|}{\Gamma_D}, \qquad \phi_{12}^{NP} = arg\left(\frac{M_{12}^{NP}}{\Gamma_{12}}\right), \tag{30}$$

Imaginary Part bounds

The LEFT (Low-Energy EFT or 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 [12] for further details).

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

$$d_u(\mu) = -2\sqrt{2}G_F \frac{eM_\tau}{\pi^2} V^2_{uD} \Im m[\hat{\epsilon}^D_T(\mu)] \log \frac{\Lambda}{\mu}\,, \tag{31}$$

  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 $\phi = \Im m\Big[V_{uD} V_{cD} \epsilon^D_{T(S)}\Big]$.

[1] We recall that [12] $d_n = g^u_T(\mu) d_u(\mu)$, 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

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

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

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

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 [12].

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 [1] the bounds from the analyses [25, 26] (see also [27, 28]), allowing the (signal-maximizing) values

$$\Re e[\epsilon_S^d] = -3.1 \times 10^{-2}\,, \quad \Re e[\epsilon_T^d] = -7.9 \times 10^{-3}\,, \tag{34}$$

and

$$\Re e[\epsilon_S^s] = 2.3 \times 10^{-2}\,, \quad \Re e[\epsilon_T^s] = 1.8 \times 10^{-2}\,. \tag{35}$$

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

In this case the SM contribution is negligible, so using current bounds for the NP part [25], we find

$$\left| A_{CP}^{rate}|_{\pi\pi} \right| \leq 3 \times 10^{-5} \,, \tag{36}$$

Clearly too tiny to be observed soon.
The maximal $A_{FB}|_{\pi\pi}$ is shown by a blue solid line in the figure.

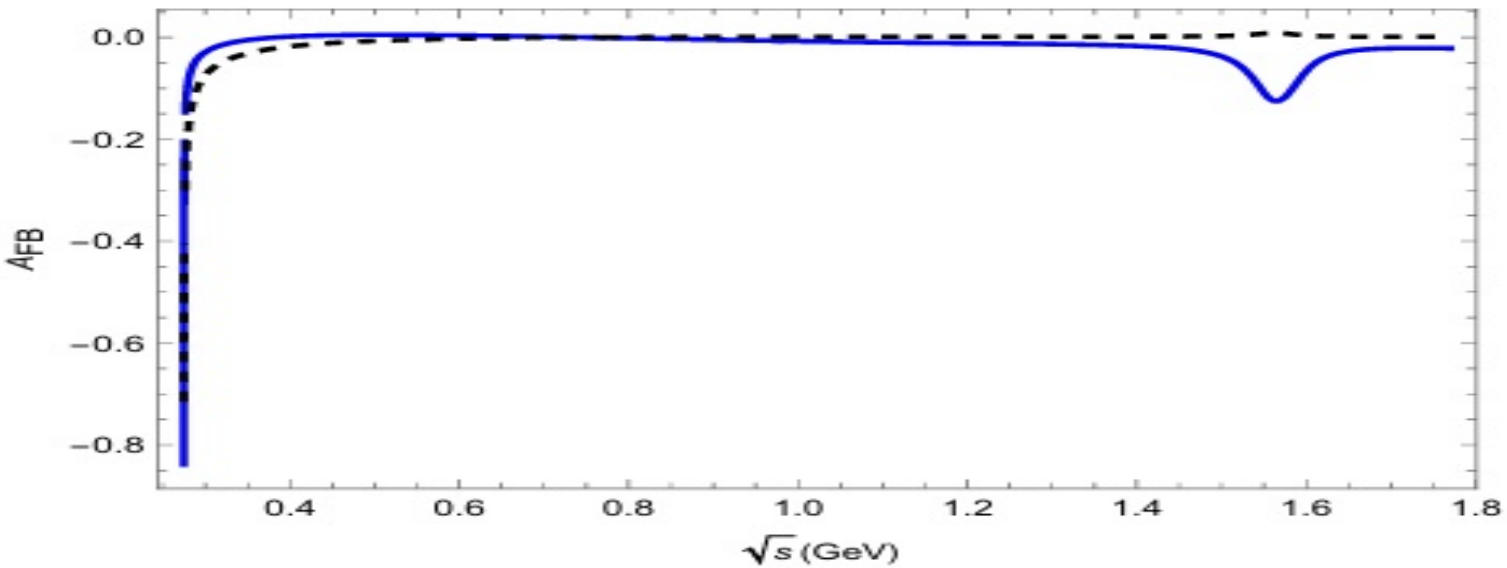

Figure: Maximal $A_{FB}^{\tau\to\pi\pi^0\nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^d] = -0.31$, $\Im m[\epsilon_S^d] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^d] = -7.9 \times 10^{-3}$ and $\Im m[\epsilon_T^d] = 8 \times 10^{-5}$), compared to the SM (black dashed line).

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

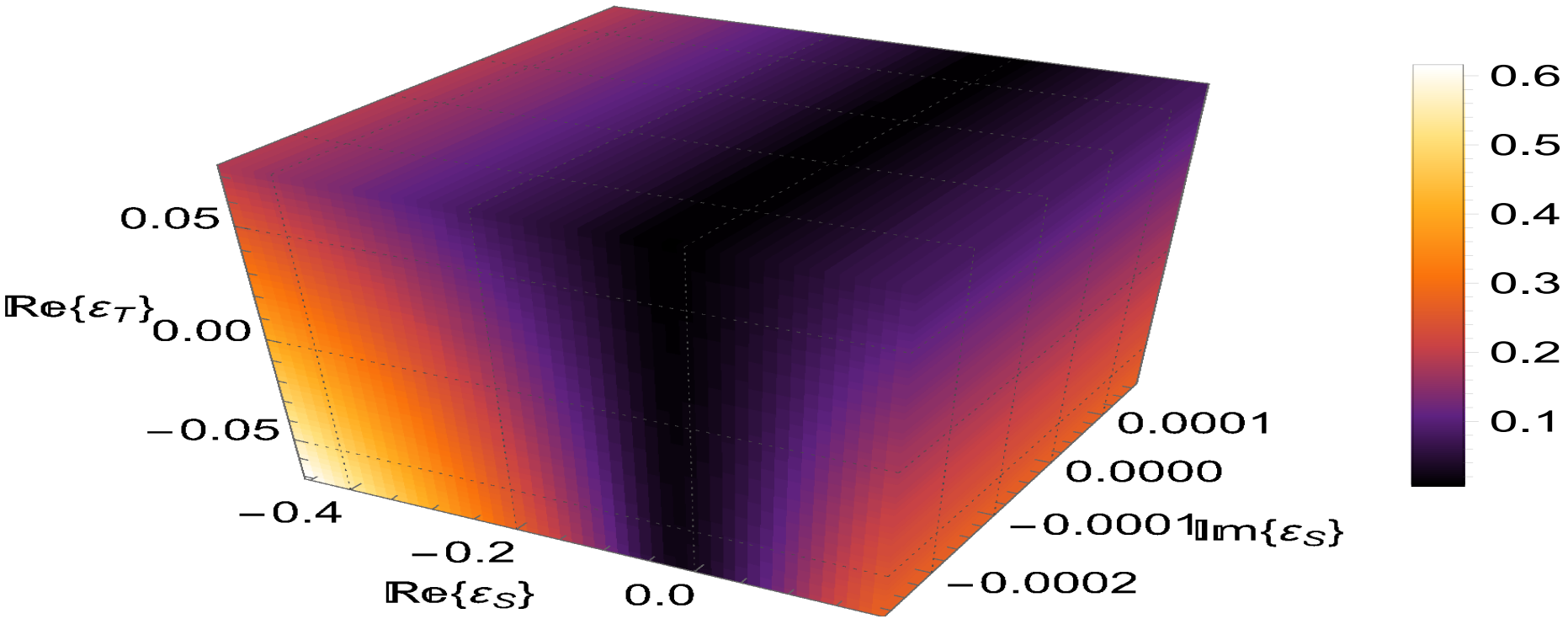

Figure: 3D density plot for the FOM on the available parameter space, setting $|\Im m[\hat{\epsilon}_T^d]| = 8 \times 10^{-5}$, for the $\pi\pi^0$ decay channel.

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 [29], whose total asymmetry rate (including $K^0 - \bar{K}^0$ oscillation effects in eq (15)) is given by

$$A^{rate}_{CP,KK} = -3.83 \times 10^{-3}. \tag{37}$$

with a maximal NP contribution given by

$$|A^{rate,NP}_{CP,KK}| \leq 2.3 \times 10^{-4} \tag{38}$$

Noticeably, a 5% precision on the measurement of A^{rate}_{CP} (feasible at FTCF) would already be sensitive to the maximum allowed NP contribution in this case.

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

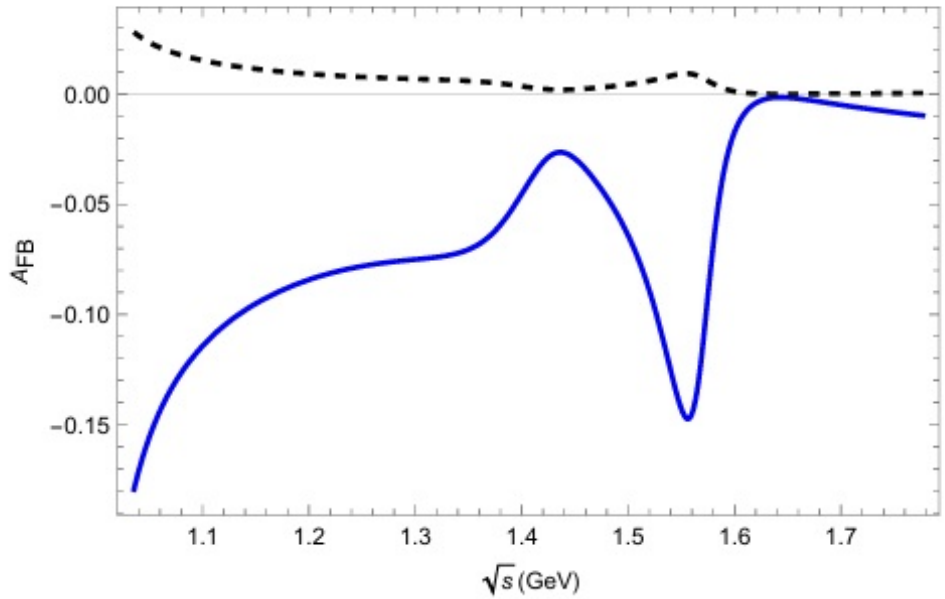

Figure: Maximal $A_{FB}^{\tau \to K_S K \nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^d] = -0.031$, $\Im m[\epsilon_S^d] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^d] = -7.9 \times 10^{-3}$ and $\Im m[\epsilon_T^d] = 8 \times 10^{-5}$, compared to the SM case (black dashed line).

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

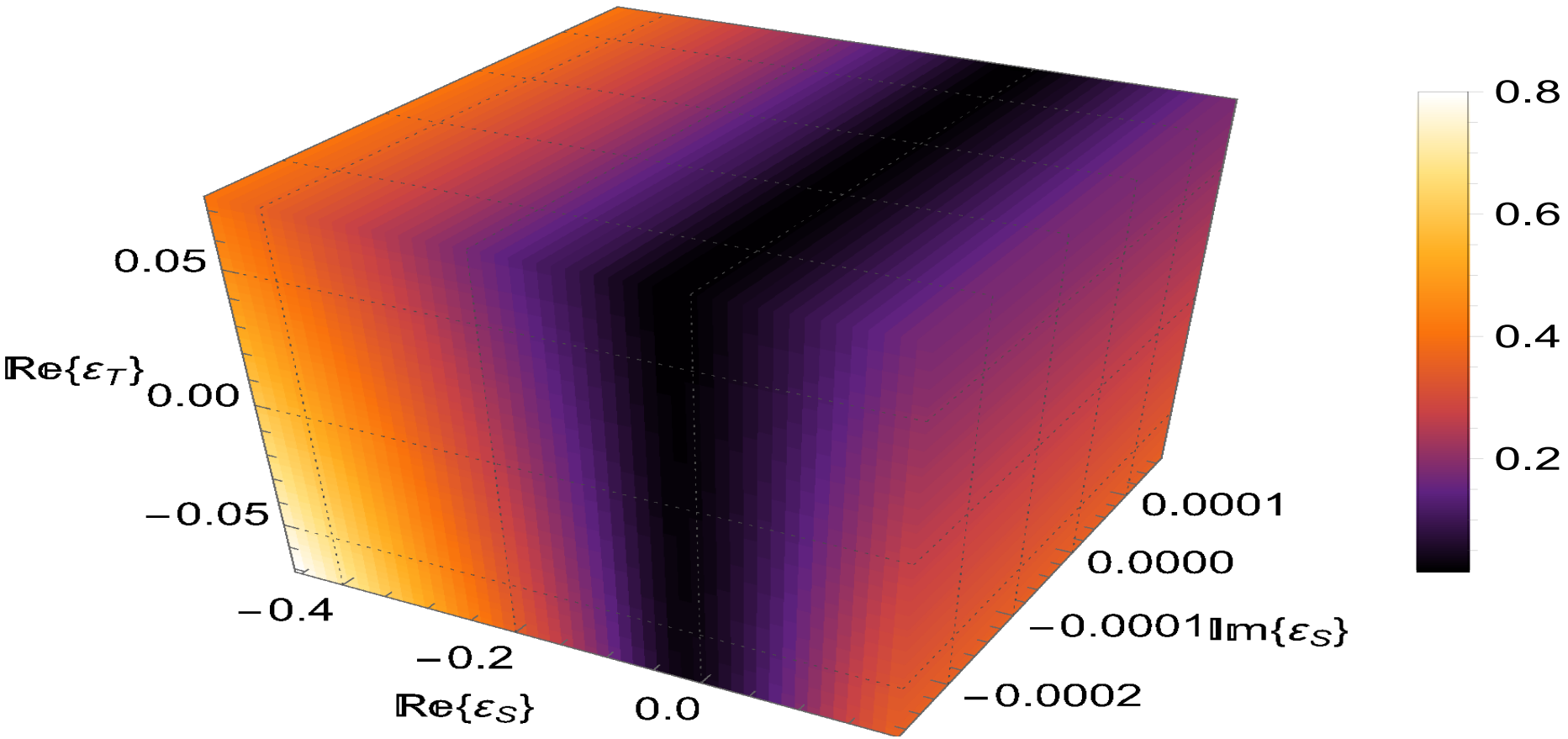

Figure: 3D density plot for the FOM described in eq. (20) on the parameter space, fixing $|\Im m[\hat{\epsilon}_T^d]| = 8 \times 10^{-5}$, for the KK_S decay channel.

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

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

$$|A_{CP}^{rate}|_{K\pi}| \leq 6 \times 10^{-7} \,. \tag{39}$$

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

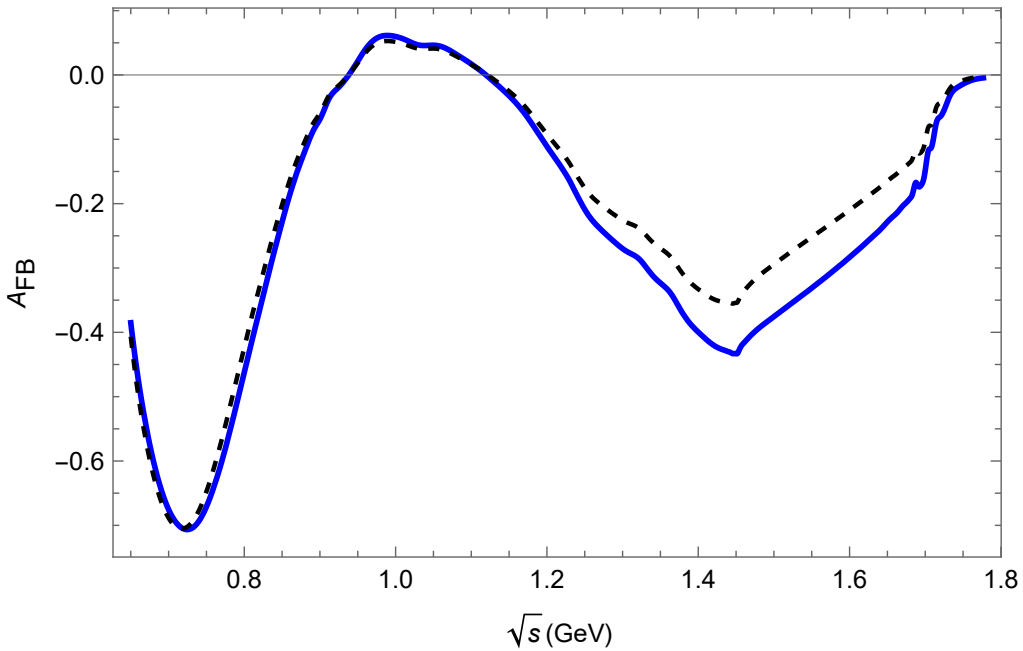

Figure: Maximal $A_{FB}^{\tau\to K\pi^0\nu_\tau}(s)$ (blue solid line), corresponding to $\Re e[\epsilon_S^s] = 2.3 \times 10^{-2}$, $\Im m[\epsilon_S^s] = -2.7 \times 10^{-4}$, $\Re e[\epsilon_T^s] = 1.8 \times 10^{-2}$ and $\Im m[\epsilon_T^s] = 4 \times 10^{-6}$ compared to the SM case (black dashed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

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

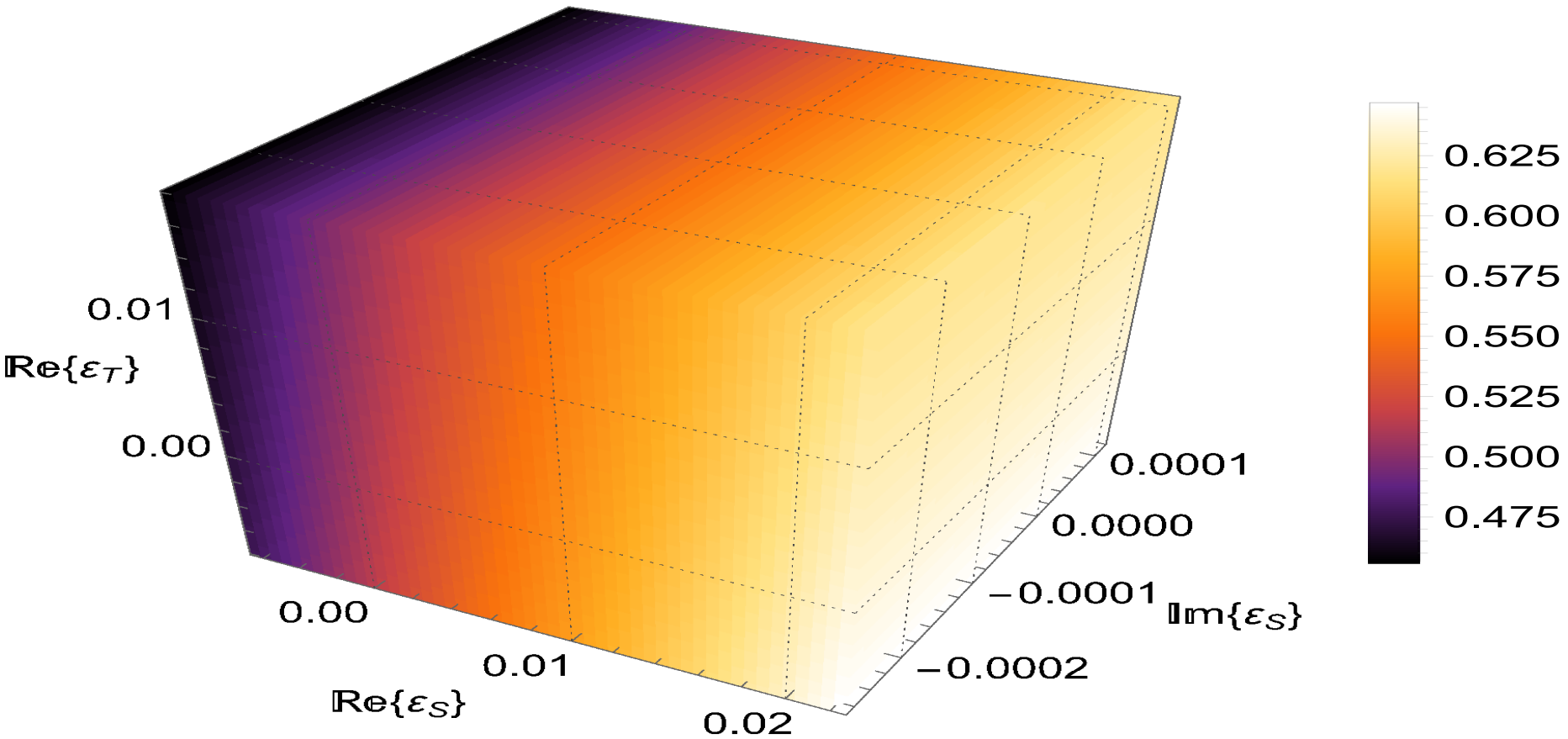

Figure: 3D density plot for the FoM described in eq. (20) on the parameter space, with $|\Im m[\hat{\epsilon}_T^s]| = 4 \times 10^{-6}$, for the decay channel with K^- and π^0 in the final state.

Conclusions

In this work we have examined (within the same framework provided by the EFT successfully applied to the $\tau \to \pi K_S \nu$ decays) 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. 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, that can reach this target accuracy. If future facilities produce polarized taus, this will open a bunch of new CP violating measurements exploiting this feature.

Supplementary material (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

$$\begin{aligned}\langle(\pi\pi)_I; out|J|\Omega\rangle &= (\langle(\pi\pi)_I; in|J|\Omega\rangle)^* \\ &= \sum_n (\langle(\pi\pi)_I; in|n; out\rangle\langle n; out|J|\Omega\rangle)^* \\ &= \langle(\pi\pi)_I; out|(\pi\pi)_I; in\rangle\,(\langle(\pi\pi)_I; out|J|\Omega\rangle)^* \\ &= e^{2i\delta_I}\,(\langle(\pi\pi)_I; out|J|\Omega\rangle)^*\,, \end{aligned} \tag{40}$$

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 scattering phase shift.

Supplementary material (Form-factors input)

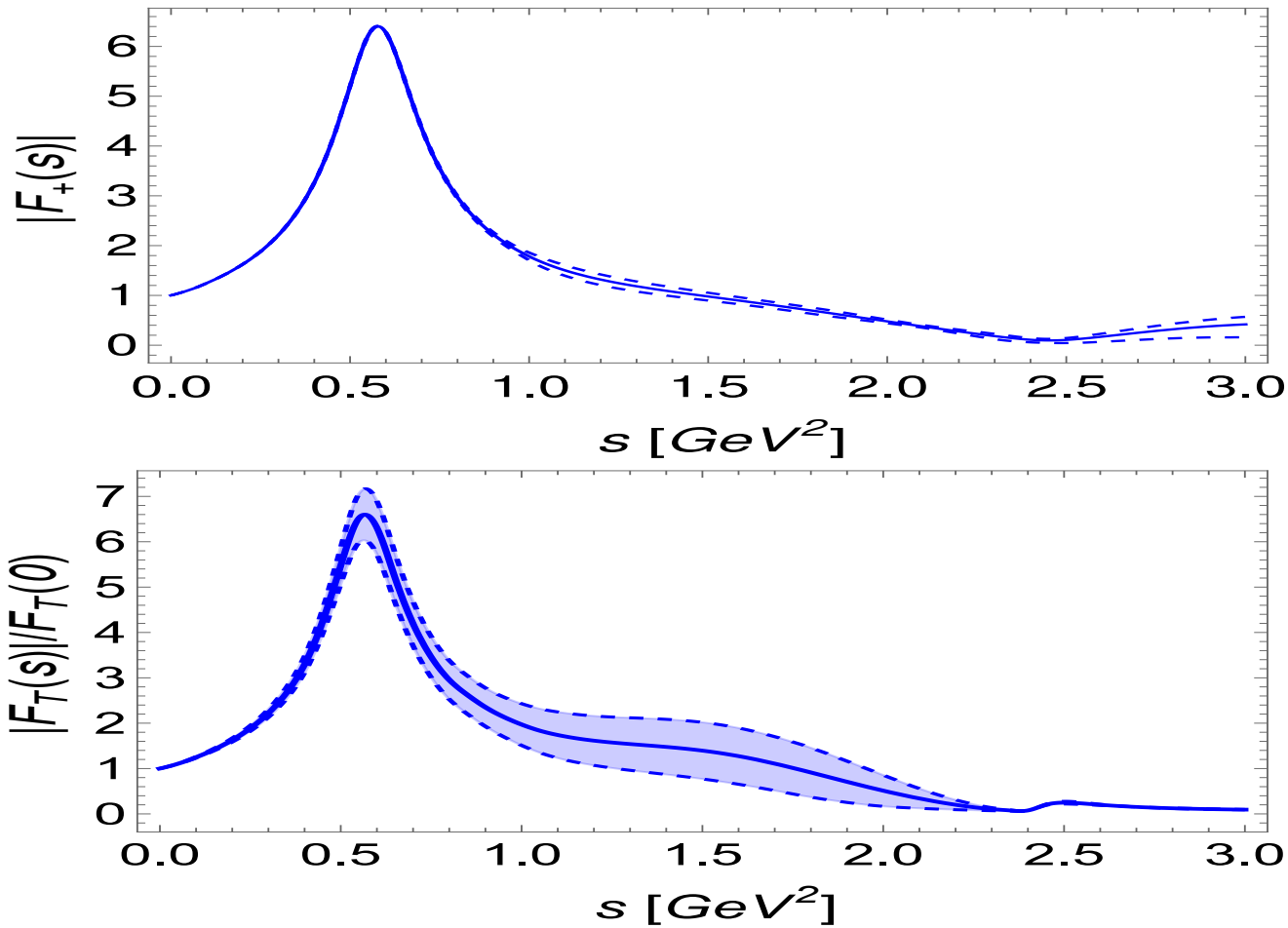

Figure: On top, pion vector form factor. At the bottom, tensor form factor derived from the estimator described in the previous section.

Supplementary material (Vector and Tensor Form-factors)

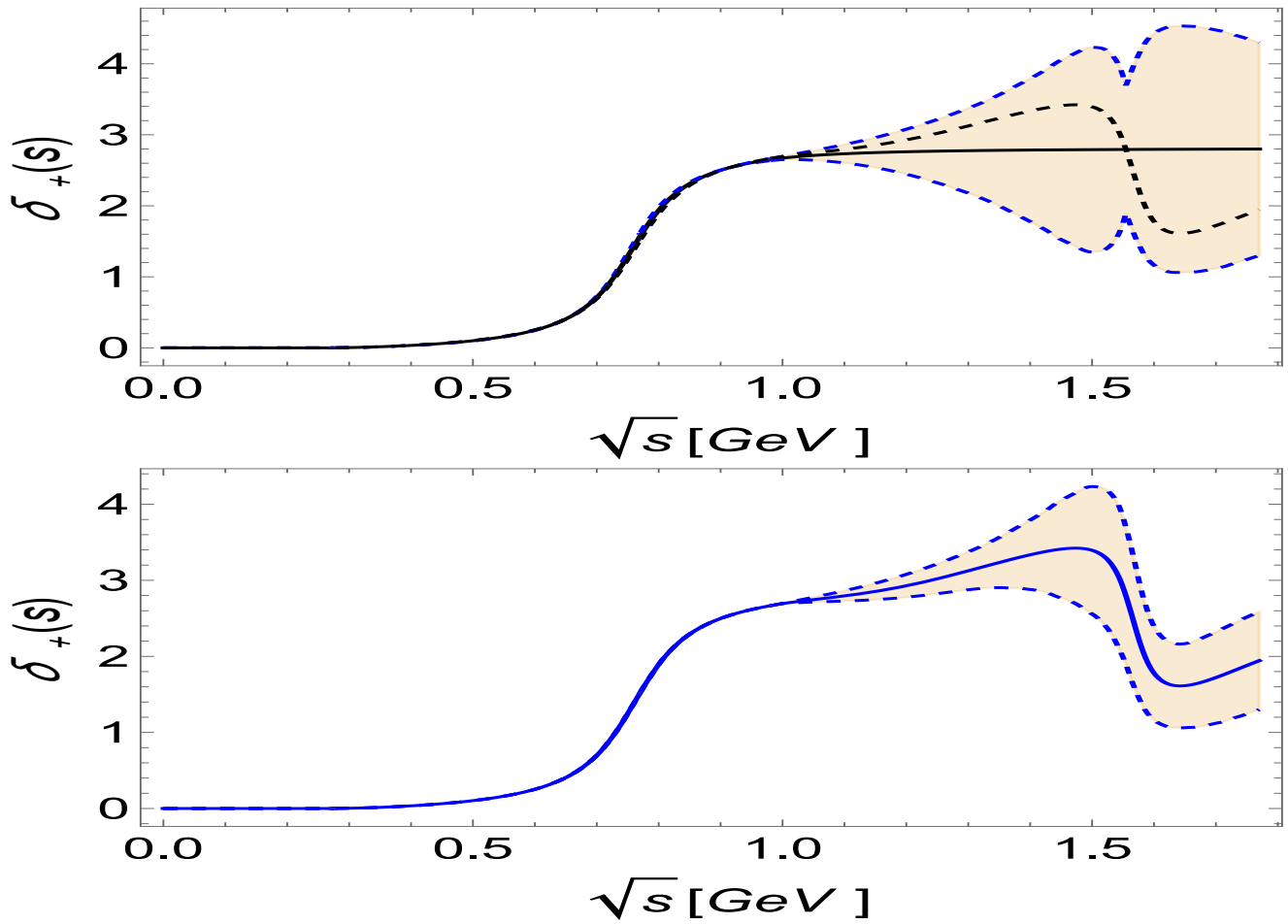

Figure: On top, tensor form factor phase. At the bottom, vector form factor phase.

- D. A. L. Aguilar, J. Rendón and P. Roig, JHEP **01** (2025), 105.
- B. Aubert *et al.* [BaBar], "Observation of CP violation in the B^0 meson system," Phys. Rev. Lett. **87** (2001), 091801.
- Y. Grossman and Y. Nir, JHEP **04** (2012), 002.
- M. Bischofberger *et al.* [Belle], Phys. Rev. Lett. **107** (2011), 131801.
- A. D. Sakharov, "Violation of CP Invariance, C asymmetry, and baryon asymmetry of the universe," Pisma Zh. Eksp. Teor. Fiz. **5** (1967), 32-35.
- H. Z. Devi, L. Dhargyal and N. Sinha, "Can the observed CP asymmetry in $\tau \to K\pi\nu_\tau$ be due to nonstandard tensor interactions?," Phys. Rev. D **90** (2014) no.1, 013016.
- V. Cirigliano, A. Crivellin and M. Hoferichter, "No-go theorem for nonstandard explanations of the $\tau \to K_S\pi\nu_\tau$ CP asymmetry," Phys. Rev. Lett. **120** (2018) no.14, 141803.

J. Rendón, P. Roig and G. Toledo Sánchez, "Effective-field theory analysis of the $\tau^- \to (K\pi)^- \nu_\tau$ decays," Phys. Rev. D **99** (2019) no.9, 093005.

F. Z. Chen, X. Q. Li, Y. D. Yang and X. Zhang, Phys. Rev. D **100** (2019) no.11, 113006.

E. A. Garcés, M. Hernández Villanueva, G. López Castro and P. Roig, JHEP **12** (2017), 027.

F. Z. Chen, X. Q. Li and Y. D. Yang, "*CP* asymmetry in the angular distribution of $\tau \to K_S \pi \nu_\tau$ decays," JHEP **05** (2020), 151.

F. Z. Chen, X. Q. Li, S. C. Peng, Y. D. Yang and H. H. Zhang, "CP asymmetry in the angular distributions of $\tau \to K_S \pi \nu_\tau$ decays. Part II. General effective field theory analysis," JHEP **01** (2022), 108.

J. H. Kühn and E. Mirkes, "Structure functions in tau decays," Z. Phys. C **56** (1992), 661-672 [erratum: Z. Phys. C **67** (1995), 364].

D. Gómez Dumm and P. Roig, "Dispersive representation of the pion vector form factor in $\tau \to \pi\pi\nu_\tau$ decays," Eur. Phys. J. C **73** (2013) no.8, 2528.

J. A. Miranda and P. Roig, "Effective-field theory analysis of the $\tau^- \to \pi^-\pi^0\nu_\tau$ decays," JHEP **11** (2018), 038.

R. Escribano, S. Gonzàlez-Solís and P. Roig, "$\tau^- \to K^-\eta^{(\prime)}\nu_\tau$ decays in Chiral Perturbation Theory with Resonances," JHEP **10** (2013), 039.

R. Escribano, S. González-Solís, M. Jamin and P. Roig, "Combined analysis of the decays $\tau^- \to K_S\pi^-\nu_\tau$ and $\tau^- \to K^-\eta\nu_\tau$," JHEP **09** (2014), 042.

S. Gonzàlez-Solís and P. Roig, "A dispersive analysis of the pion vector form factor and $\tau^- \to K^-K_S\nu_\tau$ decay," Eur. Phys. J. C **79** (2019) no.5, 436.

J. A. Miranda and P. Roig, "New τ-based evaluation of the hadronic contribution to the vacuum polarization piece of the muon anomalous magnetic moment," Phys. Rev. D **102** (2020), 114017.

G. L. Castro, A. Miranda and P. Roig, "Isospin breaking corrections in 2π production in tau decays and e+e- annihilation: Consequences for the muon g-2 and conserved vector current tests," Phys. Rev. D **111** (2025) no.7, 073004.

R. Escribano, S. Gonzàlez-Solís and P. Roig, "Predictions on the second-class current decays $\tau^- \to \pi^- \eta^{(\prime)} \nu_\tau$," Phys. Rev. D **94** (2016) no.3, 034008.

T. Bhattacharya, V. Cirigliano, R. Gupta, H. W. Lin and B. Yoon, "Neutron Electric Dipole Moment and Tensor Charges from Lattice QCD," Phys. Rev. Lett. **115** (2015) no.21, 212002.

R. Gupta, B. Yoon, T. Bhattacharya, V. Cirigliano, Y. C. Jang and H. W. Lin, "Flavor diagonal tensor charges of the nucleon from (2+1+1)-flavor lattice QCD," Phys. Rev. D **98** (2018) no.9, 091501.

C. Abel *et al.* "Measurement of the Permanent Electric Dipole Moment of the Neutron," Phys. Rev. Lett. **124** (2020) no.8, 081803.

V. Cirigliano, D. Díaz-Calderón, A. Falkowski, M. González-Alonso and A. Rodríguez-Sánchez, JHEP **04** (2022), 152.

R. Escribano, J. A. Miranda and P. Roig, "Radiative corrections to the $\tau^- \to (P_1 P_2)^- \nu_\tau \, (P_{1,2} = \pi, K)$ decays," Phys. Rev. D **109** (2024) no.5, 053003.

V. Cirigliano, A. Falkowski, M. González-Alonso and A. Rodríguez-Sánchez, "Hadronic τ Decays as New Physics Probes in the LHC Era," Phys. Rev. Lett. **122** (2019) no.22, 221801.

S. Gonzàlez-Solís, A. Miranda, J. Rendón and P. Roig, "Exclusive hadronic tau decays as probes of non-SM interactions," Phys. Lett. B **804** (2020), 135371.

S. Gonzàlez-Solís, A. Miranda, J. Rendón and P. Roig, "Effective-field theory analysis of the $\tau^- \to K^-(\eta^{(\prime)}, K^0)\nu_\tau$ decays," Phys. Rev. D **101** (2020) no.3, 034010.