Analyzing New Physics in the decays $\bar{B}^0 \to D^{(*)}\tau^-\bar{\nu}_\tau$

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Content

Experimental Status

- **Ratios of branching fractions**
  $\mathbf{R(D^{(*)}) \equiv \mathcal{B}(\bar{B}^0 \to D^{(*)}\tau^-\bar{\nu}_\tau)/\mathcal{B}(\bar{B}^0 \to D^{(*)}\mu^-\bar{\nu}_\mu)}$
- **Experiments**

$$R(D)|_{\rm BABAR} = 0.440 \pm 0.072, \qquad R(D^*)|_{\rm BABAR} = 0.332 \pm 0.030,$$
$$R(D)|_{\rm BELLE} = 0.375 \pm 0.069, \qquad R(D^*)|_{\rm BELLE} = 0.293 \pm 0.041,$$
$$R(D^*)|_{\rm LHCb} = 0.336 \pm 0.040,$$

(statistical and systematic uncertainties combined in quadrature)

- **Average ratios**

$$R(D)|_{\rm expt} = 0.391 \pm 0.050, \qquad R(D^*)|_{\rm expt} = 0.322 \pm 0.022,$$

HFAG 2015

- **SM expectations**

$$R(D)|_{\rm SM} = 0.297 \pm 0.017, \qquad R(D^*)|_{\rm SM} = 0.252 \pm 0.003,$$

Fajfer et al. 2012, Kamenik et al. 2008

$\rightarrow$ **SM excess: 1.9 σ and 3.2 σ, respectively;**

Theoretical attempts to explain the excess

1) Specific NP models: two-Higgs-doublet models (2HDMs), Minimal Supersymmetric Standard Model (MSSM), Leptoquark models, etc.

- W. S. Hou, Enhanced charged Higgs boson effects in $\mathbf{B \to \tau \bar{\nu}_\tau}$, $\mathbf{B \to \mu \bar{\nu}_\mu}$ and $\mathbf{b \to \tau \bar{\nu}_\tau + X}$, Phys. Rev. D **48**, 2342 (1993).
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- A. Crivellin, C. Greub and A. Kokulu, Explaining $\mathbf{B \to D \tau \nu}$, $\mathbf{B \to D^* \tau \nu}$ and $\mathbf{B \to \tau \nu}$ in a 2HDM of type III, Phys. Rev. D **86**, 054014 (2012).
- L. Calibbi, A. Crivellin and T. Ota, Effective field theory approach to $\mathbf{b \to s \ell \ell^{(\prime)}}$, $\mathbf{B \to K^{(*)} \nu \bar{\nu}}$ and $\mathbf{B \to D^{(*)} \tau \nu}$ with third generation couplings, Phys. Rev. Lett. **115**, 181801 (2015).
- A. Crivellin, J. Heeck and P. Stoffer, A perturbed lepton-specific two-Higgs-doublet model facing experimental hints for physics beyond the Standard Model, Phys. Rev. Lett. **116**, no. 8, 081801 (2016).
- M. Bauer and M. Neubert, One Leptoquark to Rule Them All: A Minimal Explanation for $\mathbf{R_{D^{(*)}}}$, $\mathbf{R_K}$ and $\mathbf{(g-2)_\mu}$, Phys. Rev. Lett. **116**, no. 14, 141802 (2016).
- S. Fajfer and N. Košnik, Vector leptoquark resolution of $\mathbf{R_K}$ and $\mathbf{R_{D^{(*)}}}$ puzzles, Phys. Lett. B **755**, 270 (2016).

Theoretical attempts to explain the excess

2) Model-independent approach: general SM+NP effective Hamiltonian for b $\to c\ell\nu$ is imposed

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- S. Fajfer, J. F. Kamenik, I. Nisandzic, and J. Zupan, Implications of lepton flavor universality violations in B-Decays, Phys. Rev. Lett. **109**, 161801 (2012).
- S. Fajfer, J. F. Kamenik, and I. Nisandzic, On the $\mathbf{B} \to \mathbf{D}^*\tau\bar{\nu}_\tau$ sensitivity to New Physics, Phys. Rev. D **85**, 094025 (2012).
- M. Duraisamy and A. Datta, The Full $\mathbf{B} \to \mathbf{D}^*\tau^-\bar{\nu_\tau}$ Angular Distribution and CP violating Triple Products, JHEP **1309**, 059 (2013).
- M. Duraisamy, P. Sharma and A. Datta, Azimuthal $\mathbf{B} \to \mathbf{D}^*\tau^-\bar{\nu_\tau}$ angular distribution with tensor operators, Phys. Rev. D **90**, no. 7, 074013 (2014).
- M. Tanaka and R. Watanabe, New physics in the weak interaction of $\bar{\mathbf{B}} \to \mathbf{D}^{(*)}\tau\bar{\nu}$, Phys. Rev. D **87**, 034028 (2013).
- P. Biancofiore, P. Colangelo, and F. De Fazio, On the anomalous enhancement observed in $\mathbf{B} \to \mathbf{D}^{(*)}\tau\bar{\nu}_\tau$ decays, Phys. Rev. D **87**, 074010 (2013).
- S. Bhattacharya, S. Nandi and S. K. Patra, Optimal-observable analysis of possible new physics in $\mathbf{B} \to \mathbf{D}^{(*)}\tau\nu_\tau$, Phys. Rev. D **93**, no. 3, 034011 (2016).

Effective Hamiltonian

Effective Hamiltonian for the quark-level transition $b \to c\tau^-\bar{\nu}_\tau$

$$\mathcal{H}_{eff} = 2\sqrt{2}G_F V_{cb}[(1+V_L)\mathcal{O}_{V_L} + V_R\mathcal{O}_{V_R} + S_L\mathcal{O}_{S_L} + S_R\mathcal{O}_{S_R} + T_L\mathcal{O}_{T_L}],$$

where the four-Fermi operators are written as

$$\begin{aligned}
\mathcal{O}_{V_L} &= (\bar{c}\gamma^\mu P_L b)\,(\bar{\tau}\gamma_\mu P_L \nu_\tau)\,,\\
\mathcal{O}_{V_R} &= (\bar{c}\gamma^\mu P_R b)\,(\bar{\tau}\gamma_\mu P_L \nu_\tau)\,,\\
\mathcal{O}_{S_L} &= (\bar{c}P_L b)\,(\bar{\tau}P_L \nu_\tau)\,,\\
\mathcal{O}_{S_R} &= (\bar{c}P_R b)\,(\bar{\tau}P_L \nu_\tau)\,,\\
\mathcal{O}_{T_L} &= (\bar{c}\sigma^{\mu\nu} P_L b)\,(\bar{\tau}\sigma_{\mu\nu} P_L \nu_\tau)\,.
\end{aligned}$$

Here, $\sigma_{\mu\nu} = i\,[\gamma_\mu, \gamma_\nu]\,/2$, $P_{L,R} = (1 \mp \gamma_5)/2$ are the left and right projection operators, and $V_{L,R}$, $S_{L,R}$, and T_L are the complex Wilson coefficients governing the NP contributions. In the SM one has $V_{L,R} = S_{L,R} = T_L = 0$. We assume that the neutrino is always left handed and NP only affects leptons of the third generation.

Matrix element

The invariant matrix element of $\bar{B}^0 \to D^{(*)}\tau\bar{\nu}_\tau$ can be written as

$$\begin{aligned}\mathcal{M} &= \frac{G_F V_{cb}}{\sqrt{2}}\Big[(1+V_R+V_L)\langle D^{(*)}|\bar{c}\gamma^\mu b|\bar{B}^0\rangle\bar{\tau}\gamma_\mu(1-\gamma^5)\nu_\tau \\ &+(V_R-V_L)\langle D^{(*)}|\bar{c}\gamma^\mu\gamma^5 b|\bar{B}^0\rangle\bar{\tau}\gamma_\mu(1-\gamma^5)\nu_\tau \\ &+(S_R+S_L)\langle D^{(*)}|\bar{c}b|\bar{B}^0\rangle\bar{\tau}(1-\gamma^5)\nu_\tau \\ &+(S_R-S_L)\langle D^{(*)}|\bar{c}\gamma^5 b|\bar{B}^0\rangle\bar{\tau}(1-\gamma^5)\nu_\tau \\ &+T_L\langle D^{(*)}|\bar{c}\sigma^{\mu\nu}(1-\gamma^5)b|\bar{B}^0\rangle\bar{\tau}\sigma_{\mu\nu}(1-\gamma^5)\nu_\tau\Big].\end{aligned}$$

Note that the axial and pseudoscalar hadronic matrix elements do not contribute to the $\bar{B}^0 \to D$ transition; and the scalar hadronic matrix element does not contribute to the $\bar{B}^0 \to D^*$ transition.

Form factors

Hadronic matrix elements are parametrized by a set of form factors:

$$\begin{aligned}
\langle D(p_2)|\bar{c}\gamma^\mu b|\bar{B}^0(p_1)\rangle &= F_+(q^2)P^\mu + F_-(q^2)q^\mu, \\
\langle D(p_2)|\bar{c}b|\bar{B}^0(p_1)\rangle &= (m_1+m_2)F^S(q^2), \\
\langle D(p_2)|\bar{c}\sigma^{\mu\nu}(1-\gamma^5)b|\bar{B}^0(p_1)\rangle &= \frac{iF^T(q^2)}{m_1+m_2}\left(P^\mu q^\nu - P^\nu q^\mu + i\varepsilon^{\mu\nu Pq}\right),
\end{aligned}$$

for the $\bar{B}^0 \to D$ transition, and

$$\begin{aligned}
&\langle D^*(p_2)|\bar{c}\gamma^\mu(1\mp\gamma^5)b|\bar{B}^0(p_1)\rangle \\
&= \frac{\epsilon^\dagger_{2\alpha}}{m_1+m_2}\left(\mp g^{\mu\alpha}PqA_0(q^2) \pm P^\mu P^\alpha A_+(q^2) \pm q^\mu P^\alpha A_-(q^2) + i\varepsilon^{\mu\alpha Pq}V(q^2)\right), \\
&\langle D^*(p_2)|\bar{c}\gamma^5 b|\bar{B}^0(p_1)\rangle = \epsilon^\dagger_{2\alpha}P^\alpha G^S(q^2), \\
&\langle D^*(p_2)|\bar{c}\sigma^{\mu\nu}(1-\gamma^5)b|\bar{B}^0(p_1)\rangle = -i\epsilon^\dagger_{2\alpha}\Big[\left(P^\mu g^{\nu\alpha} - P^\nu g^{\mu\alpha} + i\varepsilon^{P\mu\nu\alpha}\right)G_1^T(q^2) \\
&\quad + \left(q^\mu g^{\nu\alpha} - q^\nu g^{\mu\alpha} + i\varepsilon^{q\mu\nu\alpha}\right)G_2^T(q^2) \\
&\quad + \left(P^\mu q^\nu - P^\nu q^\mu + i\varepsilon^{Pq\mu\nu}\right)P^\alpha\frac{G_0^T(q^2)}{(m_1+m_2)^2}\Big],
\end{aligned}$$

for the $\bar{B}^0 \to D^*$ transition. Here, $P = p_1 + p_2$, $q = p_1 - p_2$, and ϵ_2 is the polarization vector of the D^* meson so that $\epsilon_2^\dagger \cdot p_2 = 0$. The particles are on their mass shells: $p_1^2 = m_1^2 = m_B^2$ and $p_2^2 = m_2^2 = m_{D^{(*)}}^2$.

Covariant Confined Quark Model in a nutshell

G. V. Efimov, M. A. Ivanov, V. E. Lyubovitskij, J. G. Körner, P. Santorelli,. . .

- **Main assumption: hadrons interact via quark exchange only**
- **Interaction Lagrangian**

$$\mathcal{L}_{int} = g_H \cdot H(x) \cdot J_H(x)$$

- **Quark current**

$$J_H(x) = \int dx_1 \int dx_2 \, F_H(x; x_1, x_2) \cdot \bar{q}^a_{f_1}(x_1) \, \Gamma_H \, q^a_{f_2}(x_2)$$

- **Vertex Function**

$$F_H(x; x_1, x_2) = \delta(x - w_1 x_1 - w_2 x_2) \, \Phi_H((x_1 - x_2)^2)$$

where $w_i = m_{q_i}/(m_{q_1} + m_{q_2})$
Translational invariant: $F_H(x + c; x_1 + c, x_2 + c) = F_H(x; x_1, x_2)$

- **Nonlocal Gaussian-type vertex functions with fall-off behavior in Euclidean space to temper high energy divergence of quark loops**

$$\widetilde{\Phi}_H(-k^2) = \int dx \, e^{ikx} \Phi_H(x^2) = e^{k^2/\Lambda_H^2}$$

where Λ_H characterizes the meson size.

Infrared confinement

General matrix element: $$\Pi = \int_0^\infty d^n\alpha\, F(\alpha_1, \ldots, \alpha_n),$$

where F stands for the whole structure of a given diagram. The set of Schwinger parameters α_i can be turned into a simplex by introducing an additional t–integration via the identity

$$1 = \int_0^\infty dt\, \delta(t - \sum_{i=1}^n \alpha_i)$$

$$\Pi = \int_0^\infty dt t^{n-1} \int_0^1 d^n\alpha\, \delta\Big(1 - \sum_{i=1}^n \alpha_i\Big)\, F(t\alpha_1, \ldots, t\alpha_n)\,.$$

Cut off the upper integration at $1/\lambda^2$

$$\Pi^c = \int_0^{1/\lambda^2} dt t^{n-1} \int_0^1 d^n\alpha\, \delta\Big(1 - \sum_{i=1}^n \alpha_i\Big)\, F(t\alpha_1, \ldots, t\alpha_n)$$

The infrared cut-off has removed all possible thresholds in the quark loop diagram.

T. Branz, A. Faessler, T. Gutsche, M. A. Ivanov, J. G. Körner and V. E. Lyubovitskij,

Model parameters

- **Model parameters are determined by fitting calculated quantities of basic processes to available experimental data or lattice simulations.**
- **The model parameters involved in this paper (all in GeV):**

$m_{u/d}$	m_s	m_c	m_b
0.241	0.428	1.67	5.04

,

λ	Λ_{D^*}	Λ_D	Λ_B
0.181	1.53	1.60	1.96

.

Form factors

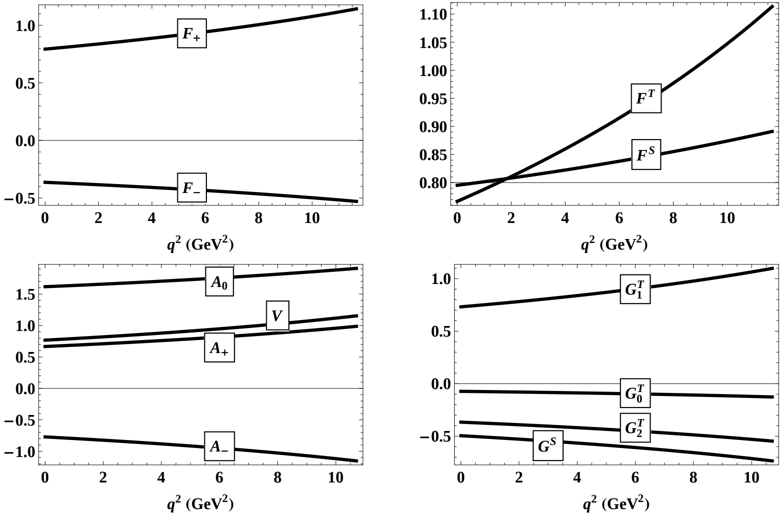

Figure : Form factors of the transitions $\bar{\mathbf{B}}^{\mathbf{0}} \to \mathbf{D}$ (upper panels) and $\bar{\mathbf{B}}^{\mathbf{0}} \to \mathbf{D}^*$ (lower panels) in the full momentum transfer range $\mathbf{0} \leq \mathbf{q}^2 \leq \mathbf{q}^2_{\mathbf{max}} = (\mathbf{m}_{\bar{\mathbf{B}}^0} - \mathbf{m}_{\mathbf{D}^{(*)}})^2$.

Form factors

- **Dipole interpolation**

$$\mathbf{F(q^2) = \frac{F(0)}{1 - as + bs^2}, \quad s = \frac{q^2}{m^2_{D^{(*)}}}.}$$

- **The dipole interpolation works very well for all form factors**

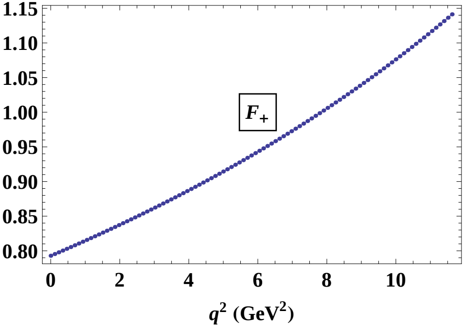

Figure : Comparison of $\mathbf{F_+(q^2)}$ form factor for $\mathbf{\bar{B}^0 \to D}$ transition calculated by FORTRAN code (dotted) with the interpolation (solid).

The parameters of the dipole interpolation:

	F_+	F_-	F^S	F^T
$F(0)$	0.79	−0.36	0.80	0.77
a	0.75	0.77	0.22	0.76
b	0.039	0.046	−0.098	0.043

,

	A_0	A_+	A_-	V	G^S	G_0^T	G_1^T	G_2^T
$F(0)$	1.62	0.67	−0.77	0.77	−0.50	−0.073	0.73	−0.37
a	0.34	0.87	0.89	0.90	0.87	1.23	0.90	0.88
b	−0.16	0.057	0.070	0.075	0.060	0.33	0.074	0.064

.

	F_+	F_-	F^S	F^T
$F(q^2_{\rm max})$	1.14	−0.53	0.89	1.11
$F^{HQL}(q^2_{\rm max})$	1.14	−0.54	0.88	1.14

,

	A_0	A_+	A_-	V	G^S	G_0^T	G_1^T	G_2^T
$F(q^2_{\rm max})$	1.91	0.99	−1.15	1.15	−0.73	−0.13	1.10	−0.55
$F^{HQL}(q^2_{\rm max})$	1.99	1.12	−1.12	1.12	−0.62	0	1.12	−0.50

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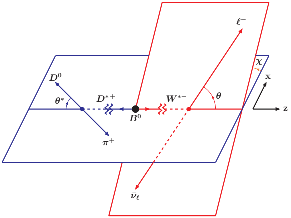
The $\bar{B}^0 \to D^{*+}(\to D^0\pi^+)\tau^-\bar{\nu}_\tau$ four-fold distribution



One has

$$\frac{d^4\Gamma(\bar{B}^0 \to D^{*+}(\to D^0\pi^+)\tau^-\bar{\nu}_\tau)}{dq^2 d\cos\theta d\chi d\cos\theta^*} = \frac{9}{8\pi}|N|^2 J(\theta,\theta^*,\chi),$$

where

$$|N|^2 = \frac{G_F^2|V_{cb}|^2|p_2|q^2v^2}{(2\pi)^3 12 m_1^2}\mathcal{B}(D^* \to D\pi).$$

The three-angle distribution

The full angular distribution $J(\theta, \theta^*, \chi)$ is written as

$$
\begin{aligned}
J(\theta, \theta^*, \chi) \\
= \quad & J_{1s} \sin^2\theta^* + J_{1c} \cos^2\theta^* + (J_{2s} \sin^2\theta^* + J_{2c} \cos^2\theta^*) \cos 2\theta \\
& + J_3 \sin^2\theta^* \sin^2\theta \cos 2\chi + J_4 \sin 2\theta^* \sin 2\theta \cos\chi \\
& + J_5 \sin 2\theta^* \sin\theta \cos\chi + (J_{6s} \sin^2\theta^* + J_{6c} \cos^2\theta^*) \cos\theta \\
& + J_7 \sin 2\theta^* \sin\theta \sin\chi + J_8 \sin 2\theta^* \sin 2\theta \sin\chi + J_9 \sin^2\theta^* \sin^2\theta \sin 2\chi,
\end{aligned}
$$

where $J_{i(a)}$ $(i = 1, \ldots, 9; a = s, c)$ are coefficient functions depending on q^2, the form factors and the NP couplings.

Experimental constraints on $\mathbf{R(D^{(*)})}$

Firstly, integrating it over all angles one obtains

$$\frac{d\Gamma(\bar{B}^0 \to D^* \tau^- \bar{\nu}_\tau)}{dq^2} = |N|^2 J_{tot} = |N|^2 (J_L + J_T),$$

where J_L and J_T are the longitudinal and transverse polarization amplitudes of the D^* meson, and given by

$$J_L = 3J_{1c} - J_{2c}, \qquad J_T = 2(3J_{1s} - J_{2s}).$$

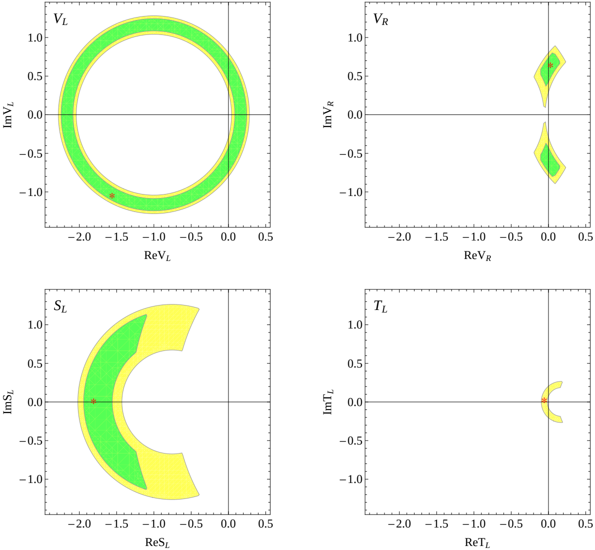
Then we calculate the ratios

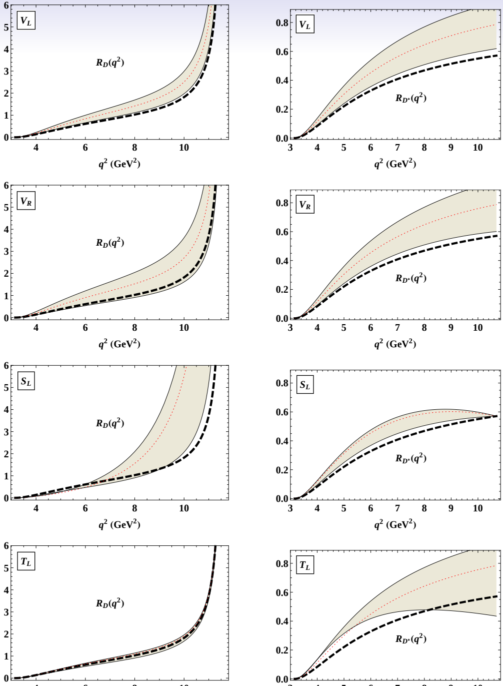
$$R_{D^{(*)}}(q^2) = \frac{d\Gamma(\bar{B}^0 \to D^{(*)} \tau^- \bar{\nu}_\tau)}{dq^2} \Big/ \frac{d\Gamma(\bar{B}^0 \to D^{(*)} \mu^- \bar{\nu}_\mu)}{dq^2}.$$

and compare with experiments to find the constraints on the space of NP couplings.

Allowed regions for NP couplings

Assuming that besides the SM contribution, only one of the NP operators is switched on at a time, and NP only affects the tau modes.





V_L
$R_D(q^2)$
q^2 (GeV^2)
V_L
$R_{D^*}(q^2)$
q^2 (GeV^2)
V_R
$R_D(q^2)$
q^2 (GeV^2)
V_R
$R_{D^*}(q^2)$
q^2 (GeV^2)
S_L
$R_D(q^2)$
q^2 (GeV^2)
S_L
$R_{D^*}(q^2)$
q^2 (GeV^2)
T_L
$R_D(q^2)$
T_L
$R_{D^*}(q^2)$

$\cos\theta$ distribution, forward-backward asymmetry & lepton-side convexity

The normalized form of the $\cos\theta$ distribution reads

$$\widetilde{J}(\theta) = \frac{a + b\cos\theta + c\cos^2\theta}{2(a + c/3)}.$$

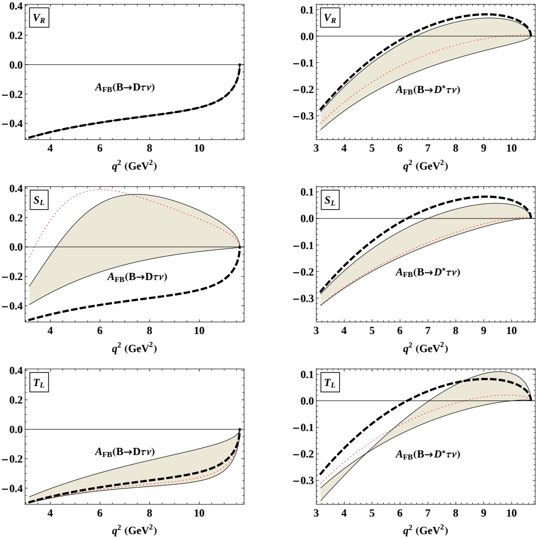
The linear coefficient $b/2(a + c/3)$ can be projected out by defining a forward-backward asymmetry given by

$$\mathcal{A}_{FB}(q^2) = \frac{(\int_0^1 - \int_{-1}^0) d\cos\theta\, d\Gamma/d\cos\theta}{(\int_0^1 + \int_{-1}^0) d\cos\theta\, d\Gamma/d\cos\theta} = \frac{b}{2(a + c/3)} = \frac{3}{2}\frac{J_{6c} + 2J_{6s}}{J_{\rm tot}},$$

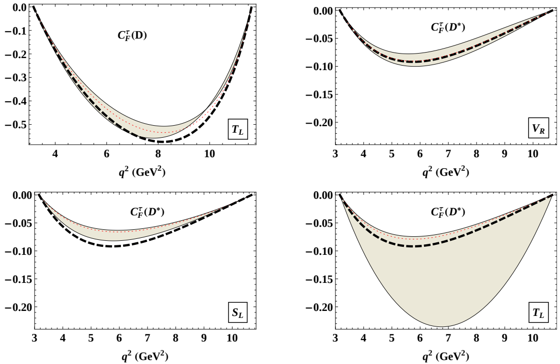
where $J_{\rm tot} = 3J_{1c} + 6J_{1s} - J_{2c} - 2J_{2s}$.
The quadratic coefficient $c/2(a + c/3)$ is obtained by taking the second derivative of $\widetilde{J}(\theta)$. Accordingly, we define a convexity parameter as follows:

$$C_F^\tau(q^2) = \frac{d^2\widetilde{J}(\theta)}{d(\cos\theta)^2} = \frac{c}{a + c/3} = \frac{6(J_{2c} + 2J_{2s})}{J_{\rm tot}}.$$

Forward-backward asymmetry $\mathcal{A}_{FB}(q^2)$



Lepton-side convexity $C_F^\tau(q^2)$



$\cos\theta^*$ distribution and hadron-side convexity parameter

The normalized form of the $\cos\theta^*$ distribution reads $\widetilde{J}(\theta^*) = (a' + c' \cos^2\theta^*)/2(a' + c'/3)$, which can again be characterized by its convexity parameter

$$C_F^h(q^2) = \frac{d^2\widetilde{J}(\theta^*)}{d(\cos\theta^*)^2} = \frac{c'}{a' + c'/3} = \frac{3J_{1c} - J_{2c} - 3J_{1s} + J_{2s}}{J_{tot}/3}.$$

The $\cos\theta^*$ distribution can be written as

$$\widetilde{J}(\theta^*) = \frac{3}{4}\left(2F_L(q^2)\cos^2\theta^* + F_T(q^2)\sin^2\theta^*\right),$$

where $F_L(q^2)$ and $F_T(q^2)$ are the polarization fractions of the D^* meson and are defined as

$$F_L(q^2) = \frac{J_L}{J_L + J_T}, \qquad F_T(q^2) = \frac{J_T}{J_L + J_T}, \qquad F_L(q^2) + F_T(q^2) = 1.$$

The hadron-side convexity parameter and the polarization fractions of the D^* meson are related by

$$C_F^h(q^2) = \frac{3}{2}\left(2F_L(q^2) - F_T(q^2)\right) = \frac{3}{2}\left(3F_L(q^2) - 1\right).$$

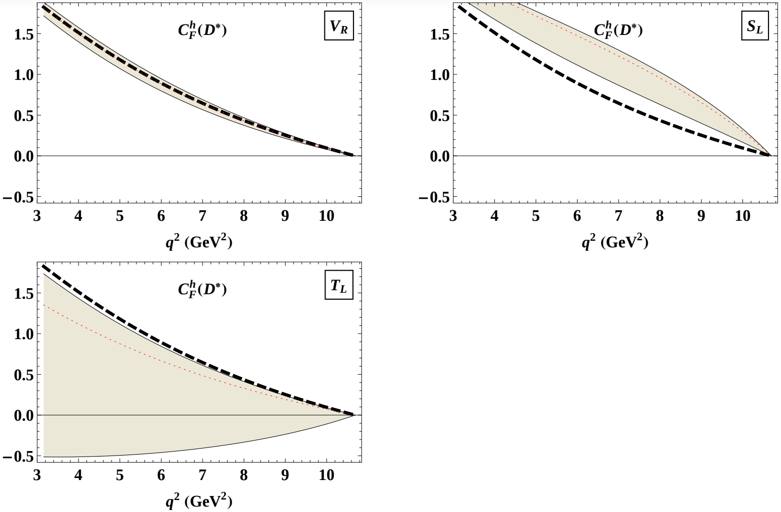


Figure : Hadron-side convexity parameter $\mathbf{C_F^h(q^2)}$.

χ distribution and trigonometric moments

The normalized χ distribution reads

$$\widetilde{J}^{(I)}(\chi) = \frac{1}{2\pi}\Big[1 + A_C^{(1)}(q^2)\cos 2\chi + A_T^{(1)}(q^2)\sin 2\chi\Big],$$

where $A_C^{(1)}(q^2) = 4J_3/J_{tot}$ and $A_T^{(1)}(q^2) = 4J_9/J_{tot}$. Besides, one can also define other angular distributions in the angular variable χ as follows

$$J^{(II)}(\chi) = \Big[\int_0^1 - \int_{-1}^0\Big] d\cos\theta^* \int_{-1}^1 d\cos\theta \frac{d^4\Gamma}{dq^2 d\cos\theta d\chi d\cos\theta^*},$$

$$J^{(III)}(\chi) = \Big[\int_0^1 - \int_{-1}^0\Big] d\cos\theta^* \Big[\int_0^1 - \int_{-1}^0\Big] d\cos\theta \frac{d^4\Gamma}{dq^2 d\cos\theta d\chi d\cos\theta^*}.$$

The normalized forms of these distributions read

$$\widetilde{J}^{(II)}(\chi) = \frac{1}{4}\Big[A_C^{(2)}(q^2)\cos\chi + A_T^{(2)}(q^2)\sin\chi\Big],$$

$$\widetilde{J}^{(III)}(\chi) = \frac{2}{3\pi}\Big[A_C^{(3)}(q^2)\cos\chi + A_T^{(3)}(q^2)\sin\chi\Big],$$

where

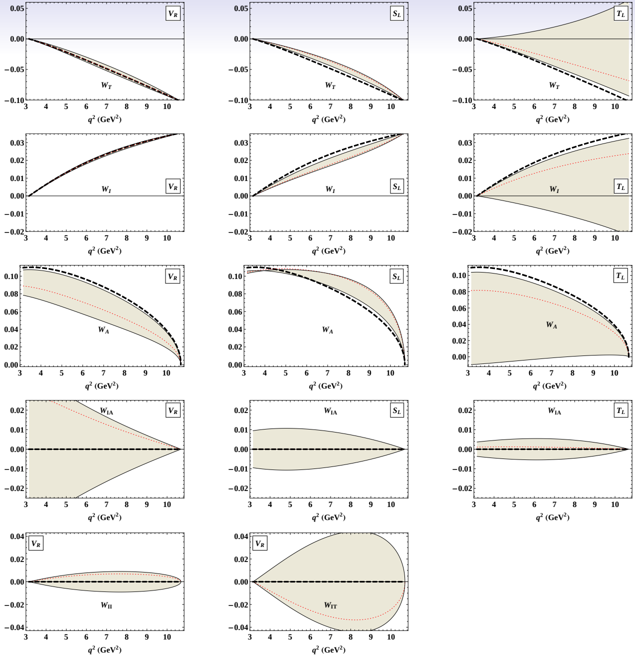
$$A_C^{(2)}(q^2) = \frac{3J_5}{J_{tot}}, \qquad A_T^{(2)}(q^2) = \frac{3J_7}{J_{tot}}, \qquad A_C^{(3)}(q^2) = \frac{3J_4}{J_{tot}}, \qquad A_T^{(3)}(q^2) = \frac{3J_8}{J_{tot}}.$$

Another method to project the coefficient functions J_i $(i = 3, 4, 5, 7, 8, 9)$ out from the full angular decay distribution is to take the appropriate trigonometric moments of the normalized decay distribution $\widetilde{J}(\theta^*, \theta, \chi)$. The trigonometric moments are defined by

$$W_i = \int d\cos\theta d\cos\theta^* d\chi M_i(\theta^*, \theta, \chi)\widetilde{J}(\theta^*, \theta, \chi) \equiv \langle M_i(\theta^*, \theta, \chi)\rangle \, ,$$

where $M_i(\theta^*, \theta, \chi)$ defines the trigonometric moment that is being taken. One finds

$$\begin{aligned}
W_T(q^2) &\equiv \langle \cos 2\chi \rangle = \frac{2J_3}{J_{\rm tot}} = \frac{1}{2}A_C^{(1)}(q^2), \\
W_{IT}(q^2) &\equiv \langle \sin 2\chi \rangle = \frac{2J_9}{J_{\rm tot}} = \frac{1}{2}A_T^{(1)}(q^2), \\
W_A(q^2) &\equiv \langle \sin\theta\cos\theta^*\cos\chi \rangle = \frac{3\pi}{8}\frac{J_5}{J_{\rm tot}} = \frac{\pi}{8}A_C^{(2)}(q^2), \\
W_{IA}(q^2) &\equiv \langle \sin\theta\cos\theta^*\sin\chi \rangle = \frac{3\pi}{8}\frac{J_7}{J_{\rm tot}} = \frac{\pi}{8}A_T^{(2)}(q^2), \\
W_I(q^2) &\equiv \langle \cos\theta\cos\theta^*\cos\chi \rangle = \frac{9\pi^2}{128}\frac{J_4}{J_{\rm tot}} = \frac{3\pi^2}{128}A_C^{(3)}(q^2), \\
W_{II}(q^2) &\equiv \langle \cos\theta\cos\theta^*\sin\chi \rangle = \frac{9\pi^2}{128}\frac{J_8}{J_{\rm tot}} = \frac{3\pi^2}{128}A_T^{(3)}(q^2).
\end{aligned}$$



V_R
S_L
T_L
W_T
W_I
W_A
W_{IA}
W_{II}
W_{IT}
q^2 (GeV^2)

Certain combinations of angular observables where the form factor dependence drops out (at least in most NP scenarios).

$$H_T^{(1)} = \frac{\sqrt{2}J_4}{\sqrt{-J_{2c}(2J_{2s} - J_3)}},$$

which equals to one not only in the SM but also in all NP scenarios except the tensor one. Therefore $H_T^{(1)}(q^2)$ plays a prominent role in confirming the appearance of the tensor operator $\mathcal{O}_{T_L}$ in the decay $\bar{B}^0 \to D^* \tau^- \bar{\nu}_\tau$.

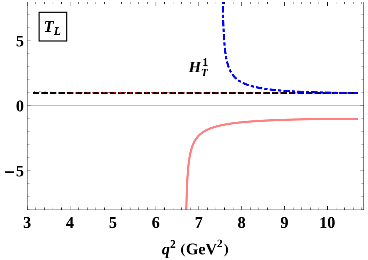


Figure : The black dashed line is the SM prediction. The red dotted line, which is almost identical to the SM one, represents the best fit value of $\mathbf{T_L}$. The blue dot-dashed line and the red line are the prediction for $\mathbf{T_L = 0.21i}$ and $\mathbf{T_L = 0.18 + 0.27i}$, respectively.

Summary and discussion

- **An analysis of possible NP in the semileptonic decays $\bar{B}^0 \to D^{(*)}\tau^-\bar{\nu}_\tau$ using the form factors obtained from our covariant quark model.**
- **Current experimental data of R(D) and R(D^*) prefer the operators $\mathcal{O}_{S_L}$ and $\mathcal{O}_{V_{L,R}}$; the operator $\mathcal{O}_{T_L}$ is less favored; and the operator $\mathcal{O}_{S_R}$ is excluded at 3σ.**
- **Our analysis can serve as a map for setting up various strategies to identify the origins of NP. For example, firstly, one uses the null-tests $W_{IT}(q^2) = 0$ and $H_T^{(1)}(q^2) - 1 = 0$ to probe the operators $\mathcal{O}_{V_R}$ and $\mathcal{O}_{T_L}$, respectively. Secondly, one measures the forward-backward asymmetry in the decay $\bar{B}^0 \to D\tau^-\bar{\nu}_\tau$. If $\mathcal{A}_{FB}^D(q^2)$ has a zero-crossing point, then it is a clear sign of the operator $\mathcal{O}_{S_L}$. The coupling V_L is more difficult to test because it is just a multiplier of the SM operator. However, if the tests above disconfirms $\mathcal{O}_{V_R}$, $\mathcal{O}_{T_L}$, and $\mathcal{O}_{S_L}$ at the same time, then the modification of V_L to R(D) and R(D^*) is a must.**

M. A. Ivanov, J. G. Körner and C. T. Tran, Exclusive decays $B \to \ell^-\bar{\nu}$ and $B \to D^{(*)}\ell^-\bar{\nu}$ in the covariant quark model, Phys. Rev. D 92, no. 11, 114022 (2015) [arXiv:1508.02678 [hep-ph]].

M. A. Ivanov, J. G. Körner and C. T. Tran, Analyzing new physics in the decays $\bar{B}^0 \to D^{(*)}\tau^-\bar{\nu}_\tau$ with form factors obtained from the covariant quark model, [arXiv:1607.02932 [hep-ph]].

Appendix: Covariant Confined Quark Model in a nutshell

G. V. Efimov, M. A. Ivanov, V. E. Lyubovitskij, J. G. Körner, P. Santorelli,...

- **Main assumption: hadrons interact via quark exchange only**
- **Interaction Lagrangian**

$$\mathcal{L}_{int} = g_H \cdot H(x) \cdot J_H(x)$$

- **Quark current**

$$J_H(x) = \int dx_1 \int dx_2\, F_H(x; x_1, x_2) \cdot \bar{q}^a_{f_1}(x_1)\, \Gamma_H\, q^a_{f_2}(x_2)$$

- **Vertex Function**

$$F_H(x; x_1, x_2) = \delta(x - w_1 x_1 - w_2 x_2)\, \Phi_H((x_1 - x_2)^2)$$

where $w_i = m_{q_i}/(m_{q_1} + m_{q_2})$

Translational invariant: $F_H(x + c; x_1 + c, x_2 + c) = F_H(x; x_1, x_2)$

- **Nonlocal Gaussian-type vertex functions with fall-off behavior in Euclidean space to temper high energy divergence of quark loops**

$$\widetilde{\Phi}_H(-k^2) = \int dx\, e^{ikx} \Phi_H(x^2) = e^{k^2/\Lambda_H^2}$$

where Λ_H characterizes the meson size.

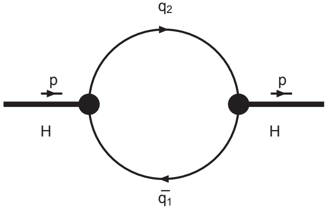
- **Compositeness condition $Z_H = 0$**

Salam 1962; Weinberg 1963

Z_H – wave function renormalization constant of the meson H.

$$Z_H^{1/2} = < H_{bare} | H_{dressed} > = 0$$

- $Z_H = 1 - \tilde{\Pi}'(m_H^2) = 0$ **where $\tilde{\Pi}(p^2)$ is the meson mass operator.**



$$\Pi_P(p) = 3g_P^2 \int \frac{dk}{(2\pi)^4 i} \widetilde{\Phi}_P^2(-k^2) \mathrm{tr}[S_1(k+w_1p)\gamma^5 S_2(k-w_2p)\gamma^5]$$

$$\Pi_V(p) = g_V^2 [g^{\mu\nu} - \frac{p^\mu p^\nu}{p^2}] \int \frac{dk}{(2\pi)^4 i} \widetilde{\Phi}_V^2(-k^2) \mathrm{tr}\,[S_1(k+w_1p)\gamma_\mu S_2(k-w_2p)\gamma_\nu]$$

The matrix elements

- **Matrix elements are described by a set of Feynman diagrams which are convolutions of quark propagators and vertex functions.**
- **Let Π be the matrix element corresponding to the Feynman diagram:**

j external momenta;
n quark propagators;
ℓ loop integrations;
m vertices.
In the momentum space it will be represented as

$$\Pi(p_1, \ldots, p_j) = \int [d^4k]^\ell \prod_{i_1=1}^{m} \Phi_{i_1+n}\left(-K^2_{i_1+n}\right) \prod_{i_3=1}^{n} S_{i_3}(\tilde{k}_{i_3} + \tilde{p}_{i_3})$$

$$K^2_{i_1+n} = \sum_{i_2} (\tilde{k}^{(i_2)}_{i_1+n} + \tilde{p}^{(i_2)}_{i_1+n})^2$$

$\tilde{k}_i$ **are linear combinations of the loop momenta** k_i

$\tilde{p}_i$ **are linear combinations of the external momenta** p_i

- **Use the Schwinger representation of the propagator:**

$$\frac{m+\not{k}}{m^2-k^2} = (m+\not{k})\int\limits_0^\infty d\alpha\ \exp[-\alpha(m^2-k^2)]$$

- **Choose a simple Gaussian form for the vertex function**

$$\Phi(-K^2) = \exp\left(K^2/\Lambda^2\right)$$

where the parameter Λ characterizes the hadron size.

- **We imply that the loop integration k proceed over Euclidean space:**

$$k^0 \to e^{i\frac{\pi}{2}} k_4 = ik_4\,, \qquad k^2 = (k^0)^2 - \vec{k}^2 \to -k_E^2 \leq 0\,.$$

- **We also put all external momenta p to Euclidean space:**

$$p^0 \to e^{i\frac{\pi}{2}} p_4 = ip_4\,, \qquad p^2 = (p^0)^2 - \vec{p}^2 \to -p_E^2 \leq 0$$

so that the quadratic momentum form in the exponent becomes negative-definite and the loop integrals are absolutely convergent.

- Convert the loop momenta in the numerator into derivatives over external momenta:

$$k_i^\mu e^{2kr} = \frac{1}{2}\frac{\partial}{\partial r_{i\,\mu}} e^{2kr},$$

- Move the derivatives outside of the loop integrals.
- Calculate the scalar loop integral:

$$\prod_{i=1}^{n}\int \frac{d^4 k_i}{i\pi^2} e^{kAk+2kr} = \prod_{i=1}^{n}\int \frac{d^4 k_i^E}{\pi^2} e^{-k_E A k_E - 2k_E r_E} = \frac{1}{|A|^2} e^{-rA^{-1}r}$$

where a symmetric $n \times n$ real matrix A is positive-definite.

- Use the identity

$$P\left(\frac{1}{2}\frac{\partial}{\partial r}\right) e^{-rA^{-1}r} = e^{-rA^{-1}r} P\left(\frac{1}{2}\frac{\partial}{\partial r} - A^{-1}r\right)$$

to move the exponent to the left.

- **Employ the commutator**

$$[\frac{\partial}{\partial r_{i\,\mu}}, r_{j\,\nu}] = \delta_{ij}\, g_{\mu\nu}$$

to make differentiation in

$$P\left(\frac{1}{2}\frac{\partial}{\partial r} - A^{-1} r\right)$$

for any polynomial P. The necessary commutations of the differential operators are done by a FORM program.

- **One obtains**

$$\Pi = \int_0^\infty d^n\alpha\, F(\alpha_1, \ldots, \alpha_n)\,,$$

where F stands for the whole structure of a given diagram.

Infrared confinement

One obtains $$\Pi = \int_0^\infty d^n\alpha \, F(\alpha_1, \ldots, \alpha_n),$$

where F stands for the whole structure of a given diagram. The set of Schwinger parameters α_i can be turned into a simplex by introducing an additional t–integration via the identity

$$1 = \int_0^\infty dt \, \delta(t - \sum_{i=1}^n \alpha_i)$$

$$\Pi = \int_0^\infty dt t^{n-1} \int_0^1 d^n\alpha \, \delta\Big(1 - \sum_{i=1}^n \alpha_i\Big) \, F(t\alpha_1, \ldots, t\alpha_n) \, .$$

Cut off the upper integration at $1/\lambda^2$

$$\Pi^c = \int_0^{1/\lambda^2} dt t^{n-1} \int_0^1 d^n\alpha \, \delta\Big(1 - \sum_{i=1}^n \alpha_i\Big) \, F(t\alpha_1, \ldots, t\alpha_n)$$

The infrared cut-off has removed all possible thresholds in the quark loop diagram.

T. Branz, A. Faessler, T. Gutsche, M. A. Ivanov, J. G. Körner and V. E. Lyubovitskij,