Lepton angular asymmetries in semileptonic charmful B decays

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Abstract

We study the lepton angular distributions in $B \to D^{(*)} \ell \nu_\ell$ decays. The lepton angular asymmetries in the decays with the general effective interactions are examined. We demonstrate that the asymmetries are sensitive to new physics with a right-handed quark current.

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Although the standard model (SM) has been accepted to be a good description of physics below the Fermi scale, it is believed that there is a more fundamental theory at a higher energy scale. Such theory will generate low energy effective couplings to various processes, which may differ from those in the SM. To find out the differences, one needs to search for physical observables which are sensitive to the new physics.

We shall concentrate on those observables in the semileptonic charmful B decays $B^- \to D^{(*)}\ell\bar{\nu}_\ell$ due to the large number of B's at the B factories. New physics effects related to the inclusive decays of $b \to c\ell\nu_\ell$ have been studied extensively in the literature [1, 2, 3, 4, 5]. In particular, the exclusive decays $B^- \to D^{(*)}\ell\bar{\nu}_\ell$ were used to constrain the scalar interactions [3] as well as the vector and axial-vector interactions [4] beyond the SM. Moreover, T violating polarization asymmetries in these exclusive modes were studied [5] in terms of all possible new interactions. In this report, we examine some asymmetrical physical observables related to the lepton angular distributions in $B^- \to D^{(*)}\ell\bar{\nu}_\ell$ with the general effective interactions.

To include the new physics effects, we start with the generalized effective Lagrangian for the process $b \to c\ell\nu_\ell$ as

$$
\begin{aligned}
\mathcal{L}_{\text{eff}} = \frac{G_F V_{cb}}{\sqrt{2}} \Big\{ & -\bar{c}\gamma_\alpha(1-\gamma_5)b\bar{\ell}\gamma^\alpha(1-\gamma_5)\nu_\ell + G_V \bar{c}\gamma_\alpha b\bar{\ell}\gamma^\alpha(1-\gamma_5)\nu_\ell \\
& +G_A \bar{c}\gamma_\alpha\gamma_5 b\bar{\ell}\gamma^\alpha(1-\gamma_5)\nu_\ell + G_S \bar{c}b\bar{\ell}(1-\gamma_5)\nu_\ell + G_P \bar{c}\gamma_5 b\bar{\ell}(1-\gamma_5)\nu_\ell + \text{h.c.} \Big\}, \quad (1)
\end{aligned}
$$

where G_F is the Fermi constant, V_{cb} is the relevant CKM matrix element, and G_i ($i = S, P, V, A$) denote the strengths of the new effective scalar, pseudoscalar, vector and axial-vector interactions, respectively. The tiny contributions from right handed neutrino are neglected. We note that, in general, the tensor interactions, described by $\bar{c}\sigma_{\mu\nu}(1 \pm \gamma_5)b\bar{\ell}\sigma^{\mu\nu}(1-\gamma_5)\nu_\ell$ with $\sigma_{\mu\nu} = i[\gamma_\mu, \gamma_\nu]/2$, should be also included in Eq. (1). However, since these tensor interactions usually arise from the models associated with baryon and lepton number violations, such as leptoqaurk models [6], for simplicity, we will not consider their effects. Instead, we will only concentrate on the models with baryon and lepton number conservations, such as the minimal supersymmetric standard model (MSSM).

To study the exclusive semileptonic B decays, such as $B^- \to D^{(*)0}\ell^-\bar{\nu}_\ell$, in terms of $\mathcal{L}$ in Eq. (1), we need to know the form factors in the transition matrix elements $\langle D^{(*)}|\bar{c}\Gamma b|B\rangle$.

We parameterize the relevant transition matrix elements to be [5],

$$\langle D(p')|\bar{c}\gamma_\mu b|B(p)\rangle \;=\; f_+(q^2)\,(p+p')_\mu + f_-(q^2)\,(p-p')_\mu\,; \tag{2}$$

$$\langle D^*(p',\epsilon)|\bar{c}\gamma^\mu b|B(p)\rangle \;=\; i\frac{F_V(q^2)}{m_B}\epsilon^{\mu\nu\alpha\beta}\epsilon_\nu^*(p+p')_\alpha q_\beta\,,$$

$$\langle D^*(p',\epsilon)|\bar{c}\gamma_\mu\gamma_5 b|B(p)\rangle \;=\; -F_{A0}(q^2)\,m_B\epsilon_\mu^* - \frac{F_{A1}(q^2)}{m_B}(p+p')_\mu\epsilon^*\cdot q - \frac{F_{A2}(q^2)}{m_B}q_\mu\epsilon^*\cdot q\,, \tag{3}$$

where p and p' are the four-momenta of B and D (D^*), respectively, ϵ is the polarization vector of D^* meson, and $q = p - p'$. By using equation of motion, the hadronic matrix elements for scalar and pseudoscalar currents are given by

$$\langle D(p')|\bar{c}b|B(p)\rangle \;=\; \frac{m_B^2}{m_b - m_c}\left[f_+(q^2)(1-r_D) + f_-(q^2)\frac{q^2}{m_B^2}\right], \tag{4}$$

$$\langle D^*(p',\epsilon)|\bar{c}\gamma_5 b|B(p)\rangle \;=\; \frac{m_B}{m_b + m_c}\epsilon^*\cdot q\left[F_{A0}(q^2) + F_{A1}(q^2)(1-r_{D^*}) + F_{A2}(q^2)\frac{q^2}{m_B^2}\right], \tag{5}$$

where $m_{b,c}$ are the quark masses and $r_{D^{(*)}} = m_{D^{(*)}}^2/m_B^2$. We note that the matrix elements $\langle D(p')|\bar{c}\gamma_\mu\gamma_5 b|B(p)\rangle$, $\langle D(p')|\bar{c}\gamma_5 b|B(p)\rangle$ and $\langle D^*(p',\epsilon)|\bar{c}b|B(p)\rangle$ are equal to zero due to parity and helicity since B (D) and D^* are pseudoscalar and vector mesons, respectively.

From the interactions in Eq. (1) and the form factors in Eqs. (2)-(5), the decay amplitudes can be written as

$$\mathcal{A}_D = \langle D\ell\nu_\ell|\mathcal{L}_{eff}|\bar{B}\rangle = \sigma\bar{\ell}(1-\gamma_5)\nu_\ell + j^\mu\bar{\ell}\gamma_\mu(1-\gamma_5)\nu_\ell \tag{6}$$

for $B \to D\ell\nu_\ell$, where $\sigma = G_S\langle D|\bar{c}b|\bar{B}\rangle$ and $j^\mu = G'_V\langle D|\bar{c}\gamma^\mu b|\bar{B}\rangle$, and

$$\mathcal{A}_{D^*} = \langle D^*(\epsilon)\ell\nu_\ell|\mathcal{L}_{eff}|\bar{B}\rangle = \Sigma(\epsilon)\bar{\ell}(1-\gamma_5)\nu_\ell + J^\mu(\epsilon)\bar{\ell}\gamma_\mu(1-\gamma_5)\nu_\ell \tag{7}$$

for $B \to D^*\ell\nu_\ell$, where $\Sigma(\epsilon) = G_P\langle D^*(\epsilon)|\bar{c}\gamma_5 b|\bar{B}\rangle$ and $J^\mu(\epsilon) = G'_V\langle D^*(\epsilon)|\bar{c}\gamma^\mu b|\bar{B}\rangle + G'_A\langle D^*(\epsilon)|\bar{c}\gamma^\mu\gamma_5 b|\bar{B}\rangle$ with $G'_{V(A)} = G_{V(A)} \mp 1$. Since our purpose is to study the lepton angular distributions, we evaluate the decay amplitudes in the rest frame of the lepton pair invariant mass q^2. The kinematical variables for particles are chosen to be

$$q = (\sqrt{q^2},0,0,0), \quad p = (E_B,0,0,|\vec{p}_X|), \quad p' = (E_X,0,0,|\vec{p}_X|),$$

$$p_\ell = (E_\ell,|\vec{p}_\ell|\sin\theta,0,|\vec{p}_\ell|\cos\theta), \quad |\vec{p}_X| = \frac{Z(m_X)}{2\sqrt{q^2}}, \quad E_B = \sqrt{|\vec{p}_X|^2 + m_B^2},$$

$$E_X = \sqrt{|\vec{p}_X|^2 + m_X^2}, \quad Z(m_X) = \sqrt{(m_B^2 - (m_X - \sqrt{q^2})^2)(m_B^2 - (m_X + \sqrt{q^2})^2)},$$

$$|\vec{p}_\ell| = (q^2 - m_\ell^2)/(2\sqrt{q^2}), \quad \epsilon(0) = \frac{1}{m_{D^*}}(|\vec{p}_{D^*}|,0,0,E_{D^*}), \quad \epsilon(\pm) = \frac{1}{\sqrt{2}}(0,1,\pm i,0), \tag{8}$$

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where X denotes D or D^*. It is clear that θ is defined as the polar angle of the lepton momentum relative to the moving direction of the B-meson in the q^2 rest frame. Hence, based on our conventions, the differential decay rate with respect to the invariant mass q^2 and the lepton polar angle for $B \to D\ell\nu_\ell$ is described by

$$
\begin{aligned}
\frac{d\Gamma}{dq^2 d\cos\theta} = {} & \frac{G_F^2 |V_{cb}|^2 |\vec{P}_{D1}||\vec{p}_\ell|^2}{2^6 \pi^3 m_B^2} \left(1 - \frac{m_\ell^2}{q^2}\right) \left\{ \frac{\sqrt{q^2}}{|\vec{p}_\ell|} \left[|\rho_D|^2 + \frac{m_\ell^2}{q^2}|G|^2 \right] + \right. \\
& \left. A_D \cos\theta + 2|G|^2 \sin^2\theta \right\} ,
\end{aligned}
\tag{9}
$$

where

$$
\begin{aligned}
A_D &= 2\frac{m_\ell}{|\vec{p}_\ell|} Re\left(\rho_D\, G^*\right), \\
\rho_D &= \sigma + \frac{m_\ell}{q^2} G'_V \left(q^2 f_-(q^2) + (m_B^2 - m_D^2) f_+(q^2)\right), \\
G &= 2|\vec{p}_D| f_+(q^2) G'_V, \quad |\vec{P}_{D1}| = \frac{Z(m_D)}{2m_B}.
\end{aligned}
\tag{10}
$$

The differential decay rate for $B \to D^*\ell\nu_\ell$ is expressed as

$$
\begin{aligned}
\frac{d\Gamma}{dq^2 d\cos\theta} = {} & \frac{G_F^2 |V_{cb}|^2 |\vec{P}_{D^*1}||\vec{p}_\ell|^2}{2^6 \pi^3 m_B^2} \left(1 - \frac{m_\ell^2}{q^2}\right) \left\{ Y_0 + A_{D^*} \cos\theta + 2|J_3|^2 \sin^2\theta \right. \\
& \left. + \left(|J_+|^2 + |J_-|^2\right) \cos^2\theta \right\} ,
\end{aligned}
\tag{11}
$$

where

$$
\begin{aligned}
A_{D^*} &= \frac{2m_\ell}{|\vec{p}_\ell|} Re\left(\rho_{D^*} J_3^*\right) + \frac{\sqrt{q^2}}{|\vec{p}_\ell|}\left(|J_+|^2 - |J_-|^2\right), \\
Y_0 &= \frac{\sqrt{q^2}}{|\vec{p}_\ell|}\left(|\rho_{D^*}|^2 + \frac{m_\ell^2}{q^2}|J_3|^2 + \frac{E_\ell}{\sqrt{q^2}}\left(|J_+|^2 + |J_-|^2\right)\right), \\
|\vec{P}_{D^*1}| &= \frac{Z(m_{D^*})}{2m_B},
\end{aligned}
\tag{12}
$$

with

$$
\begin{aligned}
\rho_{D^*} &= \Sigma + \frac{m_\ell}{\sqrt{q^2}} J_0, \\
J_0 &= -G'_A \frac{m_B}{m_{D^*}} |\vec{p}_{D^*}| \left[F_{A0}(q^2) + F_{A1}(q^2)(1 - r_{D^*}) + F_{A2}(q^2)\frac{q^2}{m_B^2} \right], \\
J_3 &= -G'_A \frac{m_B}{m_{D^*}} E_{D^*} \left[F_{A0}(q^2) + 2F_{A1}(q^2)\frac{|\vec{p}_{D^*}|^2}{m_B^2}\frac{\sqrt{q^2}}{E_{D^*}} \right], \\
J_\pm &= \frac{1}{\sqrt{2}}\left(m_B G'_A F_{A0}(q^2) \pm 2\frac{|\vec{p}_{D^*}|}{m_B}\sqrt{q^2} G'_V F_V(q^2) \right).
\end{aligned}
\tag{13}
$$

ρ_{D^*}, J_0 and J_3 denote the longitudinal contributions of D^*, while $J^\pm$ are the transverse effects.

Since the differential decay rates in Eqs. (9) and (11) involve the polar angle of the lepton, we can define an angular asymmetry to be

$$\mathcal{A}(q^2) = \frac{\int_0^{\pi/2} d\cos\theta d\Gamma/(dq^2 d\cos\theta) - \int_{\pi/2}^{\pi} d\cos\theta d\Gamma/(dq^2 d\cos\theta)}{\int_0^{\pi/2} d\cos\theta d\Gamma/(dq^2 d\cos\theta) + \int_{\pi/2}^{\pi} d\cos\theta d\Gamma/(dq^2 d\cos\theta)}, \tag{14}$$

from which we may study the behavior of the lepton angular distributions in Eqs. (9) and (11). It is easy to see that the asymmetry in Eq. (14) is related to the parity-odd terms associated with $\cos\theta = \vec{p}_X \cdot \vec{p}_\ell/|\vec{p}_X||\vec{p}_\ell|$, appearing in the formulas of the differential decay rates in Eqs. (9) and (11). Explicitly, we have

$$\mathcal{A}(q^2) \propto A_{D^{(*)}} \tag{15}$$

for $B \to D^{(*)}\ell\nu_\ell$, where $A_{D^{(*)}}$ are defined in Eqs. (10) and (12), respectively. For $B \to D\ell\bar{\nu}_\ell$ ($\ell = e, \mu, \tau$) decays, the parity-odd effects could be only generated by the interference between scalar and vector interactions. To fit the proper chirality, therefore, we need one lepton mass insertion. That is the reason why A_D in Eq. (10) is proportional to the lepton mass. Thus, the asymmetries in the electron and muon modes are negligible, but it could be large in the τ mode. However, since D^* carries the transverse degree of freedom, the parity-odd effect could be induced from such extra degree, even in the chiral limit of $m_\ell = 0$. From Eq. (3), we see that the transverse effect is related to form factors F_V and F_{A0}. One expects that A_{D^*} will be proportional to $G'_V F_V G'_A F_{A0}$. Hence, According to Eq. (12), besides the term proportional to the lepton mass, in the D^* production mode we have the contribution from $|J_+|^2 - |J_-|^2 \propto Re(G'_V G'^*_A) F_V F_{A0}$ due to the transverse polarization of D^* [8]. It is worth mentioning that if the couplings involved are symmetric in parity, i.e., $G'_V = 0$ or $G'_A = 0$, we still cannot get the angular asymmetry in $B \to D^*\ell\bar{\nu}_{\ell L}$ in the chiral limit.

To get the numerical values, the transition form factors based on the heavy quark symmetry are taken to be [5]

$$f_\pm = \pm\frac{1 \pm \sqrt{r_D}}{2r_D^{1/4}}\xi_1(w), \quad F_V = \frac{1}{2r_D^{1/4}}\xi_2(w), \quad F_{A0} = -r_D^{1/2}(1+w)\xi_2(w),$$
$$F_V = F_{A1} = -F_{A2}, \tag{16}$$

where $\xi(w)_i$ are the Isgur-Wise functions, which are normalized to unity at zero recoil, and $w = (m_B^2 + m_X^2 - q^2)/(2m_B m_X)$. We note that to include the correction due to the heavy

quark symmetry breaking, we adopt the Isgur-Wise (IW) functions for D and D^* productions as $\xi_1(w) = 1 - 0.75(w-1)$ and $\xi_2(w) = 1 - 0.95(w-1)$, respectively, based on the results in Ref. [7]. Following Eqs. (9) and (11), the decay branching ratios (BRs) for $B \to D^{(*)}\ell\bar{\nu}_\ell$ in the SM are summarized in Table I. In the table, we also show the current experimental data given by Ref. [9]. We note that the decay BRs for the light lepton modes of e and μ are insensitive to the lepton masses. The differential decay rates and angular asymmetries

TABLE I: Decay BRs (in units of 10^{-2}) of $B^- \to D^{(*)}\ell\bar{\nu}_\ell$ ($\ell = e, \mu$) and $B^- \to D^{(*)}\tau\bar{\nu}_\tau$.

Decays	$B^- \to D^0\ell\bar{\nu}_\ell$	$B^- \to D^0\tau\bar{\nu}_\tau$	$B^- \to D^{*0}\ell\bar{\nu}_\ell$	$B^- \to D^{*0}\tau\bar{\nu}_\tau$
SM	2.07	0.62	5.43	1.39
Experiments [9]	2.15 ± 0.22		6.5 ± 0.5	

in the SM are displayed in Figs. 1 and 2, respectively.
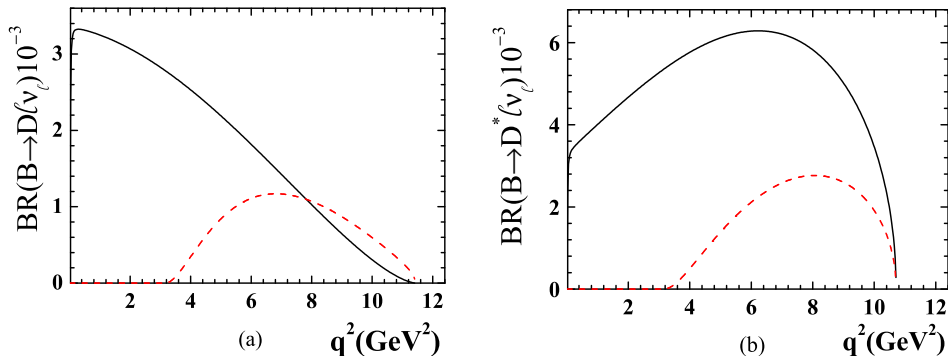

FIG. 1: Differential decay rates in the SM for (a) $B^- \to D^0\ell^-\bar{\nu}_\ell$ and (b) $B \to D^{*0}\ell^-\bar{\nu}_\ell$. The solid lines denote $\ell = e$ and μ, while the dashed lines are $\ell = \tau$.

To illustrate new physics effects on the angular asymmetries, we consider two types of new interactions. One of them is the interactions arising from a charged Higgs, described by

$$\mathcal{L}_H = \frac{G_F}{\sqrt{2}}V_{cb}C_H^\ell \bar{c}(1+\gamma_5)b\bar{\ell}(1-\gamma_5)\nu_\ell. \tag{17}$$

The other one is due to the right-handed current in the quark sector, given by

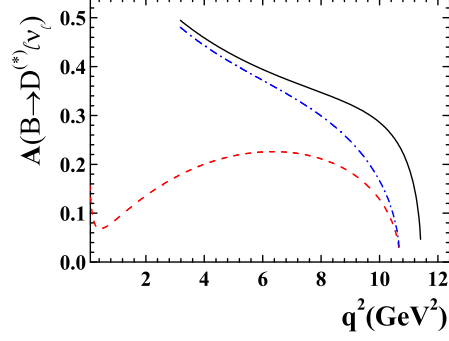$$\mathcal{L}_R = \frac{G_F}{\sqrt{2}}V_{cb}C_R\bar{c}\gamma_\mu(1+\gamma_5)b\bar{\ell}\gamma^\mu(1-\gamma_5)\nu_\ell. \tag{18}$$

FIG. 2: Angular asymmetries in the SM for $B^- \to D^0 \tau^- \bar{\nu}_\tau$ (solid), $B^- \to D^{*0} \ell^- \bar{\nu}_\ell$ ($\ell = e$ and μ, dashed) and $B^- \to D^{*0} \tau^- \bar{\nu}_\tau$ (dash-dotted).

From Eqs. (10) and (15), we find that the influence of G_S on the $B \to D \tau \bar{\nu}_\tau$ decay is much effective than that of G_V. Therefore, we only consider the contribution of $\mathcal{L}_H$ to $B \to D \tau \bar{\nu}_\tau$. On the other hand, as the transverse polarizations of D^* are sensitive to the angular asymmetries in $B \to D^* \ell \bar{\nu}_\ell$ decays, the influences of $G_{A(V)}$ are much effective than that of G_P. Hence, we only concentrate $G_{A(V)}$ for $B \to D^* \ell \bar{\nu}_\ell$ decays. It is clear that the parameters with the new interactions need to satisfy the current experimental data [9] shown in Table I. For more specific models, we adopt those governed by supersymmetry with R-parity invariance as shown in Ref. [5]. The Feynman diagrams with the tree and one-loop corrections to the SM are shown in Fig. 3. The corresponding C_H^ℓ and C_R are given
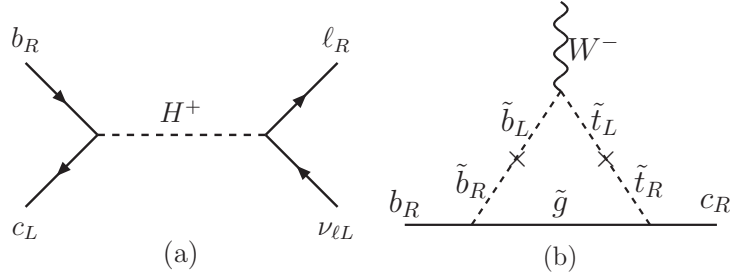

FIG. 3: Feynman diagrams for effective interactions (a) $\mathcal{L}_H$ and (b) $\mathcal{L}_R$.

by

$$C_H^\ell = -\frac{1}{m_H^2} m_b m_\ell \tan^2 \beta \,,$$

$$C_R = -\frac{\alpha_s}{36\pi} \frac{m_t(A_t - \mu \cot \beta)}{m_{\tilde{g}}^2} \frac{m_b(A_b - \mu \tan \beta)}{m_{\tilde{g}}^2} \frac{\tilde{V}_{33} V_{32}^{U_R*} V_{33}^{D_R}}{V_{cb}} I_0 \left(\frac{m_{\tilde{t}}^2}{m_{\tilde{g}}^2}, \frac{m_{\tilde{b}}^2}{m_{\tilde{g}}^2} \right) \,,$$

$$I_0(a,b) = \int_0^1 dz_1 \int_0^{1-z_1} dz_2 \frac{24 z_1 z_2}{[a z_1 + b z_2 + (1 - z_1 - z_2)]^2}, \tag{19}$$

where $A_{t(b)}$ are the soft SUSY breaking A terms, μ stands for the two Higgs superfields mixing parameter, $\tan \beta$ is the ratio of the two Higgs vacuum expectation values (VEVs), and m_i ($i = \tilde{g}, \tilde{t}, \tilde{b}, H$) are the masses of gluino, stop, sbottom and charged Higgs, respectively. Since squarks and quarks in general have different flavor structures, the unitary matrices to diagonalize the mass matrices of up(down)-quark and those of their superpartners are also different. For simplicity, we choose the bases that the mass matrices of squarks and quarks are diagonalized before including the soft SUSY breaking A terms, which govern the mixings of left handed and right-handed squarks. The effects of A terms, the 2nd and 3rd factors of C_R in Eq. (19), could be taken as perturbations. We use $\tilde{V}_{33}$ to denote the super CKM matrix associated with the coupling $W^- \tilde{b}_L^* \tilde{t}_L$ while $V^{U_R(D_R)}$ are the mixing matrices for diagonalizing the upper (down) type quarks.

Since $\tan \beta$ and the various masses are all free parameters, in the following numerical estimations, we take $\tan \beta = 50$, $m_t = 174$ GeV, $m_b = 4.4$ GeV, $m_H = 300$ GeV, $m_{\tilde{g}} = |\mu| = A_t = A_b = 200$ GeV, and $I_0 = 5$ [5, 10]; and also, to maximize C_R, we set $|\tilde{V}_{33}| = |V_{33}^{D_R}| = 1$ and $|V^{U_R}| = 1/\sqrt{2}$ [5, 10]. Hence, we obtain $C_H^\mu = -0.01$, $C_H^\tau = -0.22$ and $|C_R| \leq 0.08$. The corresponding decay BRs due to the new physics are displayed in Table II. From the

TABLE II: Decay BRs (in units of 10^{-2}) of $B^- \to D^{(*)} \ell \bar{\nu}_\ell$ with $G_S = -0.01$ and -0.22 for the μ and τ modes, respectively, and $G_{V,A} = 0.08$ (-0.08).

Decays	$B^- \to D^0 \mu \bar{\nu}_\mu$	$B^- \to D^0 \tau \bar{\nu}_\tau$	$B^- \to D^* \ell \bar{\nu}_\ell$	$B^- \to D^* \tau \bar{\nu}_\tau$
New Physics	2.07	0.84	6.25 (4.69)	1.60 (1.20)

table, we see that they are consistent with the experimental data [9]. Our results for the angular asymmetries are presented in Fig. 4. As seen from Fig. 4, the asymmetries in $B \to D^* \ell \bar{\nu}_\ell$ with $\ell = e$ and μ are more sensitive to the new physics from the right-handed

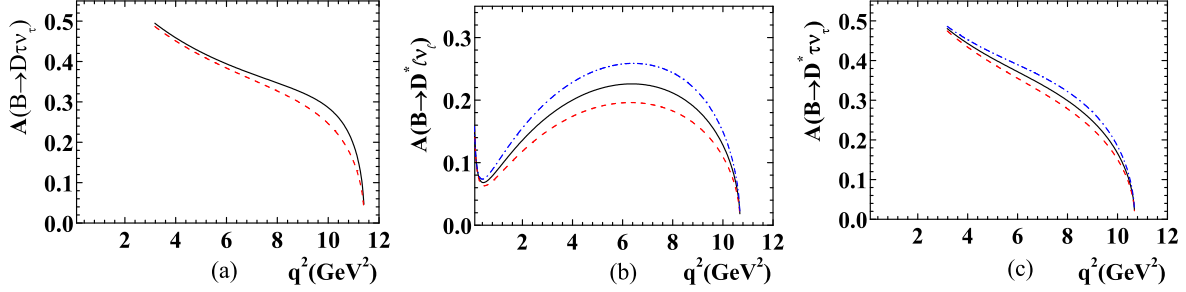FIG. 4:   Angular asymmetries with new effects for (a) $B \to D\tau\bar{\nu}_\tau$, (b) $B \to D^*\ell\bar{\nu}_\ell$ and (c) $B \to D^*\tau\bar{\nu}_\tau$. The solid lines denote the SM results. The dashed (dash-dotted) lines stand for (a) $G_S = -0.22$ and (b,c) $G_{V(A)} = 0.08$ (-0.08), respectively.

current in the quark sector. Since we adopt the IW function to be $\xi(\omega) = 1 + \rho^2(\omega - 1)$ [7, 11], by considering the experimental errors for ρ^2, we find that the shapes for angular distributions with different values of ρ^2 all overlap each other, *i.e.*, our results of Fig. 4 are insensitive to the errors from the ρ^2 parameter.

In summary, according to the lepton angular distributions in $B \to D^{(*)}\ell\nu_\ell$ decays, we have studied the angular asymmetries with the general effective interactions. We have illustrated the asymmetries in the quark currents with scalar and $V + A$ interactions, respectively. We have shown that they are sensitive to new physics with the right-handed quark current.

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