Lepton asymmetries in exclusive $b \to s \ell^+ \ell^-$ decays as a test of the Standard Model

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We argue that the longitudinal lepton polarization and the forward-backward asymmetries in exclusive $B \to (K, K^*)\ell^+ \ell^-$ decays are largely unaffected by the theoretical uncertainties in the long-distance contributions both induced by the long-range charming penguins in the $\ell^+ \ell^-$ channel and those in the meson channels encoded in the meson transition form factors, and are mainly determined by the short-distance Wilson coefficients at the low-energy scale $\mu \approx m_b$. Thus, the above mentioned asymmetries provide a powerful probe of the Standard Model and its extensions.

The investigation of rare exclusive semileptonic decays induced by the quark transition $b \to s \ell^+ \ell^-$ one faces the two types of the long-distance (LD) contributions. First, these are the LD effects in the continuum states described by the $c\bar{c}$ resonances ($\psi, \psi', ...$) and the charmed hadronic continuum states. And second, these are the LD effects in the meson transition amplitude of the Effective Hamiltonian encoded in the meson transition form factors.

In this letter we study the influence of the above mentioned two types of the LD effects on observables in exclusive $B \to K^{(*)}\ell^+ \ell^-$ decays. We show that the asymmetries of the lepton distributions remain largely unaffected by the theoretical uncertainties in the LD effects but turn out sensitive to the values of the Wilson coefficients thus providing a possibility to reliably separate and study the short-distance (SD) effects.

Neglecting the strange quark mass the effective Hamiltonian describing the $b \to s \ell^+ \ell^-$ transition has the following structure $^2$

$$
\mathcal{H}_{eff}(b \to s \ell^+ \ell^-) = \frac{G_F}{\sqrt{2}} \frac{\alpha_{em}}{2\pi} \lambda_t \left[ -2i \frac{C_7(\mu)}{q^2} m_b \bar{s}_L \sigma_{\mu\nu} q^\nu b \cdot \bar{l}_\nu l + C_{9V}(\mu) \bar{s}_L \gamma_\mu b \cdot \bar{l}_\nu l + C_{10A}(\mu) \bar{s}_L \gamma_\mu b \cdot \bar{l}_\nu l \right],
$$

where $\lambda_t = V_{tb}^* V_{ts}$ and $\bar{s}_L \equiv \bar{s}(1 + \gamma_5)$. The quantities $C_i(\mu)$ are the SD Wilson coefficients which are obtained by integrating out the heavy particles. The values of the Wilson coefficients at the scale $\mu \approx 5 \text{ GeV}$ are $^3$: $C_1(\mu) = 0.241$, $C_2(\mu) = -1.1$, $C_{7\gamma}(\mu) = 0.312$, $C_{9V}(\mu) = -4.21$ and $C_{10A}(\mu) = 4.64$.

An important contribution to the $b \to s \ell^+ \ell^-$ transition is given by the intermediate $c\bar{c}$ states in the $\ell^+ \ell^-$ channel, generated by the 4-quark operators in the effective Hamiltonian describing the $b \to s$ transition, and proceeding through the chain of transitions $b \to s c\bar{c} \to s \gamma^* \to s \ell^+ \ell^-$. The $c\bar{c}$ pairs in this channel generate both the SD and LD contributions.

The $b \to s \gamma^*$ transition vertex is transverse with respect to the photon momentum because of the gauge invariance. In the limit $m_s = 0$ within a theory with a $V - A$ charged current it has the general structure of the form

$$
\langle s(k_2) | J_{\mu}^{em}(0) | b(k_1) \rangle = \frac{2}{3} e \frac{G_F}{\sqrt{2}} \lambda_t (C_1 + C_2/3) \bar{s}_L \langle k_2 | (\gamma_\mu q^2 - q_\mu q) G_1(q^2) - 2i m_b \sigma_{\mu\nu} q^\nu G_2(q^2) \rangle b(k_1),
$$

where $G_1$ and $G_2$ are the independent form factors. It is important to point out that the form factor $G_1(q^2)$ has no pole at $q^2 = 0$. This property is related to the fact that the photon mass remains unrenormalized at all orders in perturbation theory. Also notice that the pole in $G_1(q^2)$ would have yielded a nonzero amplitude of the radiative $B \to K\gamma$ decay.
The amplitude of the $b \to s\ell^+\ell^-$ transition through the $c\bar{c}$ penguin then has the form

$$A(b \to s\ell^+\ell^-) = -\frac{G_F}{\sqrt{2}} \lambda_t (C_1 + C_2/3) \frac{2}{3} s_L \left[ \gamma_{\mu} q^2 G_1^c (q^2) - 2 i m_b \sigma_{\mu\nu} q^\nu G_2^c (q^2) \right] b \cdot \bar{l} \gamma_l \ell,$$  \hspace{1cm} (3)

where $G_{1,2}^{c\bar{c}}$ are the $c\bar{c}$ contributions to $G_{1,2}$.

The Lorentz structure of this amplitude is similar to the structure of the effective Hamiltonian \([\mathbb{H}]\), so the $c\bar{c}$ contributions to the $b \to s\ell^+\ell^-$ transition can be translated into additions to the Wilson coefficients as follows

$$C_{9V} \to C_{9V}^{c\bar{c}} (q^2) = C_{9V} + \Delta C_{9V}^{c\bar{c}} (q^2), \quad \Delta C_{9V}^{c\bar{c}} (q^2) \equiv \frac{16 \pi^2}{3} (C_1 + C_2/3) G_1^c (q^2)$$  \hspace{1cm} (4)

$$C_{7\gamma} \to C_{7\gamma}^{c\bar{c}} (q^2) = C_{7\gamma} + \Delta C_{7\gamma}^{c\bar{c}} (q^2), \quad \Delta C_{7\gamma}^{c\bar{c}} (q^2) \equiv \frac{16 \pi^2}{3} (C_1 + C_2/3) G_2^c (q^2).$$

In the $q^2$-region of the vector $c\bar{c}$ resonances ($V$) the form factors $G_{1,2}^{c\bar{c}}$ have singularities corresponding to these resonances. The general expression for the $b \to sV$ amplitude can be written as (cf. \([\mathbb{I}]\))

$$A(b \to sV) = -\frac{G_F}{\sqrt{2}} \lambda_t (C_1 + C_2/3) s_L \left[ \gamma_{\mu} M_V^2 g_1 (M_V^2) - 2 i m_b \sigma_{\mu\nu} q^\nu g_2 (M_V^2) \right] b \epsilon_\mu ^V (q),$$  \hspace{1cm} (5)

where $\epsilon_\mu ^V (q)$ is the $V$ polarization vector. Then the form factors $G_{1,2}^{c\bar{c}} (q^2)$ in the region $q^2 \approx M_V$ are expressed through the $g_{1,2}$ as follows

$$G_{1}^{c\bar{c}} (q^2) = \frac{g_1 (M_V^2)}{M_V^2 - q^2 - i \Gamma_V M_V} f_V M_V + [\text{regular terms at } q^2 = M_V^2].$$  \hspace{1cm} (6)

The leptonic decay constant $f_V$ is defined by the relation $\langle 0|\bar{c} \gamma_{\mu} c|V \rangle = \epsilon_\mu ^V (q) M_V f_V$.

If we neglect the soft nonfactorizable gluon exchanges, the amplitude of the $b \to s\ell^+\ell^-$ transition through the charming penguin is connected with the charm contribution to the vacuum polarization, $\Pi_{\mu\nu}^{c\bar{c}}$, defined as

$$\Pi_{\mu\nu}^{c\bar{c}} (q) = \frac{1}{i} \int dxe^{-iqx} \langle 0|T(c(0)\gamma_{\mu}c(0),\bar{c}(x)\gamma_{\nu}c(x))|0 \rangle.$$  \hspace{1cm} (7)

As a consequence of the vector current conservation, $\Pi_{\mu\nu}^{c\bar{c}}$ has the transverse structure

$$\Pi_{\mu\nu}^{c\bar{c}} = (g_{\mu\nu} q^2 - q_{\mu} q_{\nu}) \Pi^{c\bar{c}} (q^2).$$  \hspace{1cm} (8)

The quantity $\Pi^{c\bar{c}} (q^2)$ contains both the LD and SD contributions and does not have a pole at $q^2 = 0$ because of the gauge invariance. In this factorization approximation one has

$$G_1^{c\bar{c}} = \Pi^{c\bar{c}} (q^2), \quad G_2^{c\bar{c}} = 0.$$  \hspace{1cm} (9)

Similarly, assuming factorization for the $b \to sV$ amplitude, one finds

$$g_1 (M_V^2) = f_V / M_V, \quad g_2 (M_V^2) = 0,$$  \hspace{1cm} (10)

and relation \([\mathbb{I}]\) takes a simple form

$$G_1^{c\bar{c}} (q^2) = \frac{f_V^2}{M_V^2 - q^2 - i \Gamma_V M_V} + [\text{regular terms at } q^2 = M_V^2], \quad G_2^{c\bar{c}} (q^2) = 0,$$  \hspace{1cm} (11)

where the regular terms come from the contribution of other resonances and the $c\bar{c}$ hadronic continuum states. A standard procedure \([\mathbb{II}]\) relies in assuming the $c\bar{c}$ hadronic continuum to be described by the $c-$quark loop in the spirit of the quark-hadron duality. Then, summing over all relevant resonances and re-expressing the decay constants through the corresponding leptonic decay rates with the relation

$$\Gamma (V \to \ell^+\ell^-) = \pi \alpha_{em}^2 \frac{16}{27} \frac{f_V^2}{M_V}$$

we come to the following representation for the shift of the Wilson coefficient $C_{9V}$ in the factorization approximation (cf. \([\mathbb{III}]\))
\[ \Delta C_{9V}^{c\bar{c}}(q^2) = [3C_1(m_b) + C_2(m_b)] \cdot \left[ h(m_c/m_b, q^2/m_b^2) + \frac{3}{\kappa} \sum_{V=J/\psi, \psi'} \frac{\pi \Gamma(V \to \ell\ell) M_V}{M_V^2 - q^2 - iM_V \Gamma_V} \right]. \] (12)

The fudge factor \( \kappa \) is introduced in Eq. (12) to account for inadequacies of the naive factorization framework (see for more details). Phenomenological analyses suggest that in order to reproduce correctly the branching ratio \( \text{BR}(B \to J/\psi X \to \ell^+\ell^-X) = \text{BR}(B \to J/\psi X) \cdot \text{BR}(J/\psi \to \ell^+\ell^-) \) it should satisfy an approximate relation \( \kappa \approx 3C_1(m_b) + C_2(m_b) \approx -1 \). The SD contributions are contained in the function \( h(m_c/m_b, q^2/m_b^2) \), which describes the one-loop matrix element of the four-quark operators (see, e.g., [4] for its explicit expression).

It is known that the charm contribution described by (12) yields a strong interference between the LD and SD contributions to the decay amplitude in a broad range of \( q^2 \). It has been argued however that a simple parametrization of the hadronic \( c\bar{c} \) continuum by the quark loop in the spirit of the quark-hadron duality considerably overestimates the net effect of the charm at small \( q^2 \). One possibility to describe the \( c\bar{c} \) continuum contribution at small \( q^2 \) in a more realistic way is to take into account the regular terms in (12) by representing the LD contributions to the \( C_{9V} \) as a sum of the Breit-Wigner resonances but assuming a \( q^2 \)-dependent \( f_\psi(q^2) \). Another possibility has been implemented in Ref. [3] where the \( c\bar{c} \) contribution to \( C_{9V} \) has been described as a sum of the Breit-Wigner resonances and the continuum contribution which has been expressed through the observable cross-section of the process \( \ell^+\ell^- \to \text{charmed hadrons} \).

Since at present there are no firm arguments to determine with a sufficient accuracy the LD \( c\bar{c} \) effects we consider the difference provided by the three models of the \( \Delta C_{9V}^{c\bar{c}} \), namely, the one given by Eq. (12), and those of Refs. [3],[4], as a typical present-day theoretical uncertainty in \( C_{9V} \). The corresponding \( C_{9V}^{c\bar{c}}(q^2) \) are plotted in Fig. 1. We point out once more that the pole at \( q^2 = 0 \) in \( C_{9V}^{c\bar{c}} \) which has been discussed in [6] is ruled out by the gauge invariance.

Let us now briefly discuss the nonfactorizable effects. First, notice that factorization yields vanishing of the \( c\bar{c} \) contribution to the amplitude of the radiative \( b \to s\gamma \) decay at \( q^2 = 0 \): in this case \( G_2 = 0 \) and the structure proportional to \( G_1 \) does not contribute at \( q^2 = 0 \). Hence, the LD \( c\bar{c} \) contribution to the radiative \( b \to s\gamma \) transition is a purely nonfactorizable effect. It should be also taken into account that a nontrivial \( c\bar{c} \) contributions to the radiative \( b \to s\gamma \) decay \( \Delta C_{7V}^{c\bar{c}}(0) \neq 0 \) implies a nonzero \( \Delta C_{9V}^{c\bar{c}}(q^2) \) in the semileptonic \( b \to s\ell^+\ell^- \) decay.

A stringent test of the factorization or its breakdown is provided for instance by the polarization of \( \psi \) in the inclusive and exclusive \( b \to s\psi \) transitions: the ratio of the longitudinally-to-transversely polarized \( \psi \) does not depend on the Wilson coefficients but is rather a function of \( g_1(M_\psi^2) \) and \( g_2(M_\psi^2) \). The analysis of Ref. [3] favours the value \( |g_2(M_\psi^2)/g_1(M_\psi^2)| \approx 0.1 \pm 0.2 \), although the factorization prediction \( g_2(M_\psi^2) = 0 \) is also not ruled out. For estimating at all kinematically accessible \( q^2 \) the nonfactorizable effects which lead to a nonzero \( \Delta C_{7V}^{c\bar{c}} \) it seems reasonable to assume that the ratio \( G_2/G_1 \), which governs the size of these effects, is limited by the relation \( |G_2(q^2)/G_1(q^2)| \leq 0.2 \) in the whole kinematically accessible region of \( q^2 \). In this case, the size of the nonfactorizable effects in the region of the resonances corresponds to the analysis of [4], and at \( q^2 = 0 \) this prescription yields a nonfactorizable contribution to the \( b \to s\gamma \) decay within 5% in accordance with the estimates of Ref. [3]. In the following analysis of the sensitivity of the asymmetries in exclusive \( B \to K^{(*)}\ell^+\ell^- \) decay to the uncertainties in the LD effects we also allow for a nonfactorizable contribution \( |\Delta C_{7V}^{c\bar{c}}(q^2)| \leq 0.2|\Delta C_{9V}^{c\bar{c}}(q^2)| \).

2. The amplitude of the \( B \to K^{(*)}\ell^+\ell^- \) decay is given by the relevant mesonic matrix element of the effective Hamiltonian. The LD effects connected with the meson formation in the initial and final \( q\bar{q} \) channels are encoded in the meson transition form factors of the bilinear quark currents from the effective Hamiltonian. Various theoretical frameworks have been applied to the description of meson transition form factors: among them are constituent quark models [5],[6], QCD sum rules [8],[10], lattice QCD [12],[13], analytical constraints [14].

Lattice QCD simulations, because of its most direct connection with QCD, are expected to provide the most reliable results. However, at present the lattice calculations do not provide the form factors in the whole accessible kinematical decay region as the daughter light quark produced in \( b \) decay cannot move fast enough on the lattice and one is therefore limited to the region of not very large recoils. For obtaining form factors in the whole kinematical decay region one can use extrapolation procedures based on some parametrizations of the form factors. For instance, in [15] a simple lattice-constrained parametrization based on the constituent quark picture [12] and pole dominance is developed.

QCD sum rules give complementary information on the form factors as they can calculate the latter at not very large momentum transfers. However, in practice various versions of the QCD sum rules give remarkably different predictions, being strongly dependent on the technical subtleties of the particular version [14],[15].

Constituent quark models (QM) have proved to be a fruitful phenomenological method for the description of heavy meson transitions. A long-standing shortcoming of the quark model predictions for the form factors has been a strong dependence of the results on the QM parameters [13]. However, as we have found recently [8], a combination of the results of the lattice simulations at large \( q^2 \) and the dispersion approach based on the constituent quark picture...
allows one to considerably increase the accuracy of the predictions. Namely, constraining the parameters of the quark model by the requirement that the form factors of the QM at small recoils reproduce the results of the lattice simulations considerably decreases the uncertainty in the relevant QM parameters. Once these parameters are fixed, the spectral representations of our dispersion QM allows a direct calculation of the form factors in the whole kinematically accessible region of \(q^2\). The form factors of the dispersion QM develop the correct heavy-quark expansion at leading and next-to-leading \(1/m_Q\) orders in accordance with QCD for the transitions between heavy quarks; for the heavy-to-light transition the form factors of the dispersion QM satisfy the relations between the form factors of vector, axial-vector, and tensor currents valid at small recoil. In addition they satisfy the dispersive bounds and thus, as well as the form factor parametrizations of Ref. [18], satisfy all known rigorous theoretical constraints. However some uncertainties are still present which are connected with the errors in the results of the lattice simulations at small recoil and the approximate character of the extrapolating formulas used in [18] as well as the constituent quark picture used for the description of the form factors in Ref. [20]. We consider that the difference between the form factors of Refs. [18,20] provide a typical theoretical uncertainty of our understanding of the LD effects encoded in the transition form factors.

3. We evaluate the forward-backward \((A_{FB})\) and the longitudinal lepton polarization \((P_L)\) asymmetries for various parametrizations of the LD \(c\bar{c}\) contributions to the Wilson coefficients \(C_{9V}^{eff}\) and \(C_7\), and various sets of the form factors with formulas given in [20]. Fig. 2 shows these quantities evaluated with the GI-OGE set of the form factors from [24] and \(\Delta C_{7V} = 0\) and various prescriptions of \(C_{9V}^{eff}(q^2)\). Fig. 3 demonstrates a sensitivity of the asymmetries to the nonfactorizable effects parametrized as additions to the Wilson coefficient \(C_7\). One can see that the uncertainties due to the LD \(c\bar{c}\) effects in both \(C_7\) and \(C_{9V}\) provide a minor influence on the observable quantities in a broad kinematical region beyond the resonances. Notice that within the SM the \(P_L(B \to K^{\mp}\mu^-\mu^+) \approx 2C_{9V}C_{9V}/(C_{9V}^2 + C_{10A}^2) \simeq -1\) independently of the prescriptions chosen for the LD effects in the Wilson coefficients.

Fig. 4 presents the same asymmetries evaluated with the \(C_{9V}\) given by the Eq. (12) and \(\Delta C_{7V} = 0\) and the two sets of the form factors: namely, the lattice-constrained parametrization of Ref. [18] and the GI-OGE set from Ref. [20]. One can observe that the asymmetries turn out to be more sensitive to the variations of the transition form factors with respect to the uncertainties in the LD effects induced by the \(c\bar{c}\) penguins. Nevertheless, there is still a broad range of momentum transfers in which the interference between the LD and SD contributions is negligible. Thus, a study of the asymmetries under consideration in these regions provides a potential possibility to measure the Wilson coefficients.

In fact, a sensitivity of the asymmetries to the SD contributions encoded in the Wilson coefficients might be observed. The CLEO data on the radiative exclusive and inclusive \(b \to s\gamma\) decays allow \(R_{7\gamma} = C_{7\gamma}(M_W)/C_{7\gamma}(M_W)\) in the ranges (see, e.g. [22]) \(0.4 \leq R_{7\gamma} \leq 1.2\), and \(-4.2 \leq R_{7\gamma} \leq -2.4\). As discussed in Refs. [22,23] scanning over the MSSM parameter space correspondent to the allowed \(C_{7\gamma}\) regions provides only minor changes in \(C_{9V}(M_W)\) and \(C_{10A}(M_W)\). Thus for illustrating possible new physics effects in lepton asymmetries we fix \(C_{9V}(M_W)\) and \(C_{10A}(M_W)\) to their SM values, and vary \(C_{7\gamma}\) in the allowed regions. The corresponding \(A_{FB}\) and \(P_L\) are shown in Fig. 5. One finds the shape of \(A_{FB}\) to be sensitive to the \(C_{7\gamma}\) value (or at least to the \(C_{7\gamma}\) sign), and this effect far overwhelms the uncertainties in the LD contributions.

A weaker but still visible sensitivity of \(P_L(B \to K^{\ast}e^\pm\nu^\mp)\) to the SD contributions can be seen (Fig. 5. b).

The longitudinal lepton polarization asymmetry \(P_L(B \to K^{\ast}\mu^\pm\nu^\mp)\) is equal to \(P_L \simeq 2C_{9V}C_{10A}/(C_{9V}^2 + C_{10A}^2)\) if \(C_{7\gamma}\) lies in the allowed regions at all kinematically accessible \(q^2\), except for the end-points and regions near \(\psi\) and \(\psi^\prime\). Therefore, the measurement of \(P_L(B \to K^{\ast}\mu^\pm\nu^\mp)\) can provide us direct information on the ratio \(C_{9V}/C_{10A}\).

In conclusion, we have analyzed the sensitivity of the lepton asymmetries in rare exclusive \(b \to s\ell^+\ell^-\) transitions to the LD effects. We have found the most important uncertainty to come from the errors in the meson transition form factors whereas the uncertainties in the LD effects induced by the long-range \(c\bar{c}\) penguins are of relatively minor importance. So for obtaining better predictions more accurate calculations of the transition form factors in the whole accessible \(q^2\) region are necessary. Nevertheless, there are broad \(q^2\)-intervals where the interference between the LD and SD effects is negligible and the asymmetries in these regions contain information on the SD effects.

In particular, the shape of the forward-backward asymmetry in \(B \to K^{\ast}\mu^\pm\nu^\mp\) turns out to be specific for the particular values of the SD Wilson coefficients: namely, within the SM, \(A_{FB}\) is positive at small \(q^2\), has a zero at \(q^2 \approx 0.15M_B^2\) and then becomes negative at larger \(q^2\). At the same time, the present experimental restrictions allow a region of the \(C_{7\gamma}\) with a sign opposite to the SD value which yields an essentially different shape of the \(A_{FB}\) (cf. also [24]). The longitudinal lepton polarization asymmetry \(P_L(B \to K^{\ast}\mu^\pm\nu^\mp)\) also contains important information on the Wilson coefficients at the scale \(\mu \simeq m_b\), and \(P_L(B \to K^{\ast}\mu^\pm\nu^\mp)\) directly measures \(C_{9V}/C_{10A}\) at \(\mu \simeq m_b\).

Thus, the experimental study of the forward-backward asymmetry and the longitudinal lepton polarization asymmetry potentially provides an effective test of the Standard Model and its possible extensions.

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FIG. 1. The effective Wilson coefficients $C_{\text{eff}}(q^2)$ for various parametrizations of the $c\bar{c}$ contributions: solid - Eq. (12), dotted - [7], dashed - [8]. The bold solid line corresponds to $C_{\text{eff}}$ without the LD contributions, namely, $C_{\text{eff}}(\mu) = (3C_1 + C_2)\ln(m_c/m_b) + \ln(q^2/m_b^2)$, the horizontal line corresponds to $-C_{\text{eff}}(\mu = 5 \text{ Gev}) = 4.21$. 

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FIG. 2. A sensitivity of the $A_{FB}$ in $B \rightarrow K^* \mu^+ \mu^-$ (a), $P_L$ in $B \rightarrow K^* \ell^+ \ell^-$ (b), and $P_L$ in $B \rightarrow K \ell^+ \ell^-$ (c) to the uncertainty in $C_{9V}^{ij}$. The GI-OGE Set of the form factors $[20]$ and $\Delta C_{77}^{CS} = 0$ are used. The curves correspond to various parametrizations of $C_{9V}^{ij}(q^2)$. Notations see Fig. 1.
FIG. 3. A sensitivity of the $A_{FB}$ in $B \to K^*\mu^+\mu^-$ (a), $P_L$ in $B \to K^*\ell^+\ell^-$ (b), and $P_L$ in $B \to K\ell^+\ell^-$ (c) to the nonfactorizable effects parametrized as additions to $C_{7\gamma}$: The $C_{9}^{f\ell j}$ from [8] and the GI-OGE set of the form factors [20] are used. The curves correspond to $\Delta C_{7\gamma}^{c\bar{c}} = 0$ (dashed) $\Delta C_{7\gamma}^{c\bar{c}} = 0$ (solid) $\Delta C_{7\gamma}^{c\bar{c}} = -0.2\Delta C_{9Y}^{c\bar{c}}$ (dotted). Bold line corresponds to $C_{9Y}$ without the LD contributions and $\Delta C_{7\gamma}^{c\bar{c}} = 0$. 
FIG. 4. A sensitivity of the $A_{FB}$ (a) and $P_L$ (b) in $B \to K^* \ell^+ \ell^-$ to the meson transition form factors: solid - GI-OGE Set from Ref. [20]; dotted - lattice-constrained parametrizations from Ref. [18]. The $C_{9V}^{eff}$ from [8] and $\Delta C_{7g}^{eff}$ are used. Bold lines correspond to the $C_{9V}$ without LD contributions.
FIG. 5. A sensitivity of the $A_{FB}$ (a) and $P_L$ (b) in $B \to K^\ast \mu^+\mu^-$ to the Wilson coefficient $C_{7\gamma}$: the asymmetries are evaluated for the GI-OGE Set of the form factors [20] and different values of $R_{7\gamma} = C_{7\gamma}(M_W)/C_{7\gamma}^{SM}(M_W)$ from the allowed region [22]: upper dotted line - $R_{7\gamma} = 1.2$, lower dotted line - $R_{7\gamma} = 0.4$, upper dashed line - $R_{7\gamma} = -2.4$, lower dashed line - $R_{7\gamma} = -4.2$, solid line - $R_{7\gamma} = 1$ (SM).