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Forward-backward and isospin asymmetry for $B \to K^*\ell^+\ell^-$ decay in the standard model and in supersymmetry

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Abstract: We discuss two dedicated observables in exclusive $B \to K^*\ell^+\ell^-$ decay that can be used to study effects of physics beyond the standard model, namely the forward-backward asymmetry in the lepton spectrum and the isospin-asymmetry between decays of charged and neutral B mesons. We consider the region of large recoil-energy (i.e. small invariant mass of the lepton pair), and employ the QCD factorization approach to exclusive B meson decays. Sub-leading effects in the heavy quark mass expansion are taken into account for the calculation of the isospin-asymmetry. We give predictions for decay asymmetries in the standard model, and its supersymmetric extension with minimal flavor violation, using parameter values allowed by current experimental constraints on $B \to X_s\gamma$ decay.

Keywords: B-Physics, Rare Decays, NLO Computations.

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1. Introduction

B physics has become one of the most active fields in the search for new interactions associated to physics beyond the standard model (SM). In particular, the flavour sector of the new theory should have distinctive effects on exclusive and inclusive rare B decays which are suppressed in the SM. At present, one of the most important constraints on new physics from rare B decays arises from the inclusive decay $B \to X_s\gamma$.

The inclusive decay processes $B \to X_s\gamma$ and also $B \to X_s l^+l^-$ have been studied in great detail in the literature in the SM [1]–[6] and MSSM [7]–[16]. Concerning $B \to X_s\gamma$ we have at our disposal accurate determinations both from experimental [17] and theoretical point of view [1]–[5]. Within the present uncertainties $B \to X_s\gamma$ is consistent with the SM prediction, which translates into a quite narrow allowed region for the Wilson coefficient $C_7^{\rm eff}$ in the weak effective hamiltonian for $b \to s\gamma$ (also $b \to s\gamma\gamma$ [18]) and $b \to s\ell^+\ell^-$ transitions. The inclusive $B \to X_s l^+l^-$ decay has not been measured yet; for the theoretical literature we refer the reader to [6] and [13]–[16].

In this paper we will focus on the exclusive decay mode $B \to K^*l^+l^-$. In general, exclusive decays are much easier accessible in experiment (first experimental evidence for this decay mode has been presented in [19, 20]), but their theoretical treatment is usually more difficult because of non-perturbative hadronic binding effects. However, in the limit where the invariant mass of the lepton pair is small, the process admits a systematic theoretical description using QCD factorization in the heavy quark limit [21]. As stressed in [22]–[25], a distinct observable in this decay is the forward-backward (FB) asymmetry in the lepton spectrum: In particular the FB asymmetry-zero can be calculated with rather high accuracy since most of the hadronic uncertainties drop out. This leads to an almost model-independent relation between the Wilson coefficients $C_7^{\rm eff}$ and C_9.

We further stress that the sensitivity to new physics in exclusive modes may be different from that of their inclusive counterparts if the spectator quark is involved in the short-distance dynamics. This is the case for the isospin asymmetry between $B \to K^* \ell^+ \ell^-$ decays with charged or neutral B mesons. Kagan and Neubert [26] have shown that annihilation topologies provide the main source for the isospin asymmetry in $B \to K^* \gamma$ decays, and concluded that the sign and magnitude of the isospin asymmetry is predominantly determined by the Wilson coefficients C_5 and C_6 of four-quark penguin operators and their relative sign with $C_7^{\rm eff}$. The SM prediction is consistent with experiment, within still large uncertainties. In this article we will extend the analysis of [26] to the decay $B \to K^* \ell^+ \ell^-$ and calculate the isospin-asymmetry for small invariant lepton-pair masses (q^2 below $7\,{\rm GeV}^2$).

The paper is organized in the following way. In the next section 2 we will shortly review the FB asymmetry for the decay $B \to K^* l^+ l^-$ within the QCD factorization approach. Then we turn to the calculation of the isospin asymmetry in $B \to K^* l^+ l^-$, and give a numerical estimate within the SM. In section 3 we explore new physics contributions to both observables, the FB and the isospin-asymmetry, in the MSSM with minimal flavor violation (MFV). We apply the constraints from $B \to X_s \gamma$ using the NLO calculation in [10, 11]. Section 4 is devoted to a brief discussion of possible effects beyond MFV, before we conclude with a summary in section 5.

2. $B \to K^* \ell^+ \ell^-$ decay asymmetries in the SM

The decay $B \to K^* \ell^+ \ell^-$ is induced by a set of operators $\mathcal{O}_i$ appearing in the weak effective hamiltonian for $b \to s$ transitions [4]

$$H_{\rm eff} = \frac{G_F}{\sqrt{2}} \left\{ \sum_{i=1}^{2} (\lambda_u \, C_i \, O_i^u + \lambda_c \, C_i \, O_i^c) - \lambda_t \sum_{i=3}^{10} C_i \, O_i \right\} . \tag{2.1}$$

Here, the combinations of CKM matrix elements are abbreviated as $\lambda_q = V_{qs}^* \, V_{qb}$, and obey the unitarity relation $\lambda_u + \lambda_c = -\lambda_t$. The Wilson coefficients C_i contain the short-distance physics related to integrating out heavy particles (like weak gauge bosons and the top quark) and hard gluons. In extensions of the SM new (heavy) particles modify the SM values of Wilson coefficients and/or lead to new operators in addition to (2.1). In this paper we will restrict ourselves to the case of SM operators (2.1) only. (We will use the same conventions as in [22]; in particular, the gauge vertices read $+iee_q$ and $+ig_s T^A$, and the SM values for $C_7^{\rm eff}$ and $C_8^{\rm eff}$ are negative.)

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A systematic calculation of QCD effects in exclusive B decays requires the factorization of long- and short-distance physics distinguished by the heavy quark mass m_b in the hadronic matrix elements of $H_{\rm eff}$. The theoretical framework has been introduced for non-leptonic B decays by Beneke/Buchalla/Neubert/Sachrajda (BBNS) [21], and has been extended to heavy-to-light form factors and radiative B decays in [22] and [26]–[31]. In the BBNS approach, starting from the heavy quark limit, $m_b \to \infty$, radiative corrections to "naive" factorization of hadronic amplitudes arising from hard gluon exchange can be calculated in terms of perturbative hard-scattering kernels and universal (non-perturbative) hadronic quantities (decay constants, form factors, and light-cone distribution amplitudes).

In order to constrain new-physics scenarios in exclusive decays with reasonable accuracy, one has to find observables that are sensitive to the short-distance physics and rather independent to hadronic input parameters like form factor values, moments of light-cone wave functions etc. In this paper we will focus on two asymmetries in the lepton spectrum, namely the forward-backward asymmetry

$$\frac{dA_{\rm FB}}{dq^2} \equiv \frac{1}{d\Gamma/dq^2} \left(\int_0^1 d(\cos\theta)\, \frac{d^2\Gamma[B \to K^*\ell^+\ell^-]}{dq^2 d\cos\theta} - \int_{-1}^0 d(\cos\theta)\, \frac{d^2\Gamma[B \to K^*\ell^+\ell^-]}{dq^2 d\cos\theta} \right), \tag{2.2}$$

and the (CP averaged) isospin asymmetry

$$\frac{dA_I}{dq^2} \equiv \frac{d\Gamma[B^0 \to K^{*0}\ell^+\ell^-]/dq^2 - d\Gamma[B^\pm \to K^{*\pm}\ell^+\ell^-]/dq^2}{d\Gamma[B^0 \to K^{*0}\ell^+\ell^-]/dq^2 + d\Gamma[B^\pm \to K^{*\pm}\ell^+\ell^-]/dq^2}\,. \tag{2.3}$$

Because of the sensitivity to relative signs and orders of magnitude between different Wilson coefficients, these asymmetries are particularly useful to constrain new physics scenarios. ([32] entertains another possibility to study physics beyond the SM in this B decay mode, by comparing helicity amplitudes for $B \to K^*\ell^+\ell^-$ and $B \to \rho\ell^+\nu$.)

2.1 Forward-backward asymmetry

We will follow the results and notation of [22] where the QCD corrections to the decay $B \to K^*\ell^+\ell^-$ have been calculated to leading power in the inverse heavy quark mass or light-quark energy, respectively, and to NLO in the strong coupling constant. The approach is restricted to the situation where the energy of K^* is large, i.e. the invariant lepton-pair mass q^2 is small. For practical purposes this means that q^2 is below charm threshold, $q^2 < 4m_c^2 \approx 7\,{\rm GeV}^2$. The decay amplitude for $B \to K^*\ell^+\ell^-$ is conveniently expressed in terms of the quantities $\mathcal{C}_9^\perp(q^2)$ and $\mathcal{C}_9^\parallel(q^2)$, which are generalizations of the effective Wilson coefficient $C_9^{\rm eff}$ used in the "naive" factorization approach. The lowest-order (naively factorizing) expressions read

$$\begin{aligned} \mathcal{C}_9^{(0)\,\perp}(q^2) &= C_9 + Y(q^2) + \frac{2m_b M_B}{q^2}\, C_7^{\rm eff}\,, \\ \mathcal{C}_9^{(0)\,\parallel}(q^2) &= C_9 + Y(q^2) + \frac{2m_b}{M_B}\, C_7^{\rm eff}\,. \end{aligned} \tag{2.4}$$

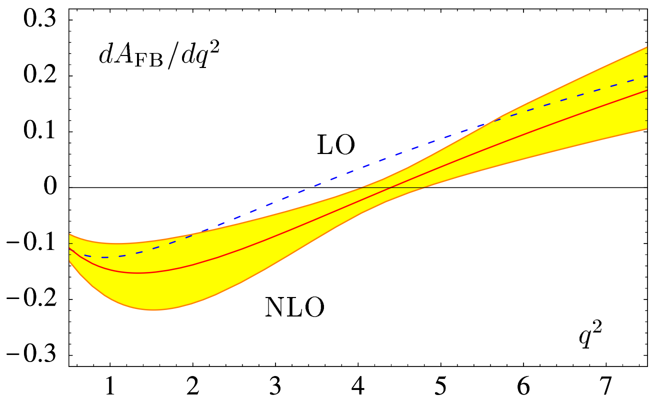

Figure 1: The SM forward-backward asymmetry $dA_{\rm FB}/dq^2$ for the decay $B^- \to K^{*-}\ell^+\ell^-$ as a function of the invariant lepton-pair mass q^2 at next-to-leading order (solid center line) and leading order (dashed). The band reflects all theoretical uncertainties from input parameters and scale dependence combined. Figure taken from [22].

(The function $Y(s)$ can be found, for instance, in [4, 22]). The differential decay rate then reads [22]

$$\begin{aligned}\frac{d^2\Gamma}{dq^2 d\cos\theta} &= \frac{G_F^2|V_{ts}^*V_{tb}|^2}{128\pi^3}\, M_B^3\,\lambda(q^2,m_{K^*}^2)^3\left(\frac{\alpha_{\rm em}}{4\pi}\right)^2 \times \\ &\times\left[(1+\cos^2\theta)\,\frac{2q^2}{M_B^2}\,\xi_\perp(q^2)^2\left(|\mathcal{C}_9^\perp(q^2)|^2+C_{10}^2\right)+\right.\\ &\quad+(1-\cos^2\theta)\left(\frac{E\,\xi_{||}(q^2)}{m_{K^*}}\right)^2\left(|\mathcal{C}_9^{||}(q^2)|^2+C_{10}^2\,\Delta_{||}(q^2)^2\right)-\\ &\quad\left.-\cos\theta\,\frac{8q^2}{M_B^2}\,\xi_\perp(q^2)^2\,\mathrm{Re}(\mathcal{C}_9^\perp(q^2))\,C_{10}\right]. \end{aligned} \tag{2.5}$$

Here

$$\lambda(q^2,m_{K^*}^2) = \left[\left(1-\frac{q^2}{M_B^2}\right)^2 - \frac{2m_{K^*}^2}{M_B^2}\left(1+\frac{q^2}{M_B^2}\right)+\frac{m_{K^*}^4}{M_B^4}\right]^{1/2} \tag{2.6}$$

is the usual kinematic factor, $m_{b,\rm PS}$ is the potential-subtracted heavy quark mass [33], $E \simeq (M_B^2-q^2)/2M_B$ is the K^* energy, and θ refers to the angle between the positively charged lepton and the B meson in the center-of-mass frame of the lepton pair. Furthermore $\xi_\perp(q^2)$ and $\xi_{||}(q^2)$ are the universal soft heavy-to-light form factors, arising in the heavy-quark and large-recoil-energy limit [34, 27] (for a critical examination whether the form factors themselves can be calculated perturbatively, see [35] and references therein.) The function $\Delta_{||}(q^2)$ is given in [22].

From (2.5) it is obvious that the FB asymmetry $dA_{\rm FB}/dq^2$ in (2.2) is proportional to $\mathrm{Re}\left(\mathcal{C}_9^\perp(q^2)\right)$, and therefore it vanishes if $\mathrm{Re}\left(\mathcal{C}_9^\perp(q_0^2)\right)=0$. At leading order this gives a relation between the short-distance Wilson coefficients C_9 and $C_7^{\rm eff}$ which is completely

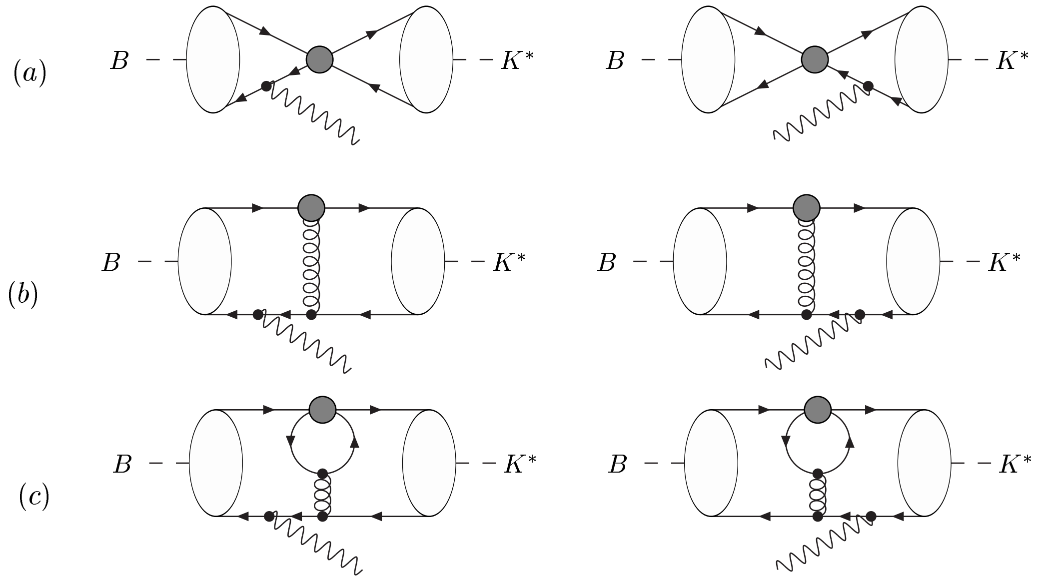

Figure 2: Contributions to the isospin asymmetry dA_I/dq^2: (a) Annihilation topologies with operators $\mathcal{O}_{1-6}$, (b) Hard spectator interaction involving the gluonic penguin operator $\mathcal{O}_8$, (c) Hard spectator interaction involving the operators $\mathcal{O}_{1-6}$.

independent of hadronic uncertainties [23, 24]

$$\text{LO:} \qquad C_9 + \mathrm{Re}(Y(q_0^2)) = -\frac{2M_B m_b}{q_0^2}\, C_7^{\mathrm{eff}}\,. \tag{2.7}$$

Therefore the measurement of the asymmetry zero provides a crucial test of the SM and new physics. (Phenomenological consequences for the FB asymmetry from physics beyond the SM have already been discussed in [24] using "naive" factorization.) In figure 1 we plot the results for the FB asymmetry obtained in [22] for SM Wilson coefficients. A conservative estimate for the location of the asymmetry zero in the SM (taking into account the possibility of additional model-dependence from the form factors arising at higher order in the Λ_h/m_B expansion) was obtained as

$$q_0^2 = (4.2 \pm 0.6)\,\mathrm{GeV}^2\,. \tag{2.8}$$

2.2 Isospin asymmetry

The isospin-asymmetry arises from "non-factorizable" graphs where a photon is radiated from the spectator quark in annihilation or spectator-scattering diagrams, see figure 2. The corresponding contribution to the partial decay width depends on the charge of the spectator quark, and therefore it is different for charged and neutral B meson decays. The isospin asymmetry vanishes in naive factorization. Non-factorizable annihilation and spectator-scattering topologies induce only a tiny asymmetry when working in the heavy quark limit. However, it has been observed in [26] that for $B \to K^*\gamma$ decays, isospin-breaking effects, that formally are of sub-leading order in the Λ_h/m_B expansion, numerically may be large.

Although part of the calculation requires to model new non-perturbative effects, the dominant isospin-breaking effect stems from processes which are dominated by short-distance effects and are thus perturbatively calculable in the QCD factorization approach.

In the following, we extend the analysis in [26] to the $B \to K^*\ell^+\ell^-$ case, including also longitudinal photon polarizations appearing for $q^2 \neq 0$. For that purpose we consider the exclusive quantities $\mathcal{C}_9^\perp$ and $\mathcal{C}_9^{||}$ introduced in [22], and define

$$\mathcal{C}_9^\perp(q^2) = \left[1 + b_q^\perp(q^2)\right] \mathcal{C}_9^{(0)\,\perp}(q^2)\,, \qquad \mathcal{C}_9^{||}(q^2) = \left[1 + b_q^{||}(q^2)\right] \mathcal{C}_9^{(0)\,||}(q^2)\,, \tag{2.9}$$

where the leading-order expressions $\mathcal{C}_9^{(0)\,\perp}(q^2)$ and $\mathcal{C}_9^{(0)\,||}(q^2)$ are given in (2.4), and the functions $b_q^\perp(q^2)$ and $b_q^{||}(q^2)$ parameterize non-factorizing effects from photon radiation off the spectator quark line with flavor q. Non-factorizable effects that are independent of the spectator quark will drop out of the isospin asymmetry to the considered order of the α_s and $1/m_b$ expansion. (Note that there is a formally LO annihilation contribution to $\mathcal{C}_9^{||}$ which does depend on the charge e_q of the spectator quark and will therefore appear in the function $b_q^{||}(q^2)$. Because this contribution is numerically small we will treat it as a sub-leading effect.) In terms of these new parameters the (CP-averaged) isospin asymmetry (2.3) can be written as

$$\begin{aligned}\frac{dA_I}{dq^2} \simeq\ & \mathrm{Re}(b_d^\perp - b_u^\perp)\ \frac{|\mathcal{C}_9^{(0)\,\perp}|^2}{|\mathcal{C}_9^{(0)\,\perp}|^2 + |\mathcal{C}_{10}|^2} \times \\ & \times \left(1 + \frac{1}{4}\cdot\frac{E^2 m_B^2}{q^2 m_{K^*}^2}\cdot\frac{\xi_{||}^2(q^2)}{\xi_\perp^2(q^2)}\cdot\frac{\mathrm{Re}(b_d^{||} - b_u^{||})}{\mathrm{Re}(b_d^\perp - b_u^\perp)}\cdot\frac{|\mathcal{C}_9^{(0)\,||}|^2}{|\mathcal{C}_9^{(0)\,\perp}|^2}\right) \times \\ & \times \left(1 + \frac{1}{4}\cdot\frac{E^2 m_B^2}{q^2 m_{K^*}^2}\cdot\frac{\xi_{||}^2(q^2)}{\xi_\perp^2(q^2)}\cdot\frac{|\mathcal{C}_9^{(0)\,||}|^2 + |\mathcal{C}_{10}|^2}{|\mathcal{C}_9^{(0)\,\perp}|^2 + |\mathcal{C}_{10}|^2}\right)^{-1}\,,\end{aligned} \tag{2.10}$$

where we neglected terms of order b_q^2. In the limit $q^2 \to 0$ the photon-pole in $\mathcal{C}_9^{(0)\,\perp}$ will dominate, $\mathcal{C}_9^{(0)\,\perp} \sim 1/q^2$, see (2.4). In this limit the isospin-asymmetry in (2.10) reduces to $A_I[B \to K^*\gamma] = \mathrm{Re}\left[b_d^\perp(0) - b_u^\perp(0)\right]$, which is the quantity discussed in [26]. For larger values of q^2 the isospin asymmetry will be dominated by the longitudinal polarization.

Generalizing the definitions in [26], we further define the functions $K_{1,2}^\perp$ and $K_1^{||}$ via

$$\begin{aligned} b_q^\perp(q^2) &= \frac{24\,\pi^2\, m_B\, f_B\, e_q}{q^2\,\xi_\perp(q^2)\,\mathcal{C}_9^{(0)\,\perp}(q^2)} \left\{\frac{f_{K^*}^\perp}{m_B}\, K_1^\perp(q^2) + \frac{f_{K^*} m_{K^*}}{6\lambda_{B,+}(q^2) m_B}\,\frac{K_2^\perp(q^2)}{1 - q^2/m_B^2}\right\}, \\ b_q^{||}(q^2) &= \frac{24\,\pi^2\, f_B\, e_q\, m_K^*}{m_B\, E\, \xi_{||}(q^2)\,\mathcal{C}_9^{(0)\,||}(q^2)} \left\{\frac{f_{K^*}}{3\,\lambda_{B,-}(q^2)}\, K_1^{||}(q^2)\right\}. \end{aligned} \tag{2.11}$$

Here

$$\lambda_{B,\pm}^{-1}(q^2) = \int_0^\infty d\omega\, \frac{\phi_\pm^B(\omega)}{\omega - q^2/m_B - i\epsilon}\,, \tag{2.12}$$

are moments of the B meson light-cone distribution amplitudes as defined for instance in [36, 27]. The function $K_1^{||}$ receives an annihilation contribution to leading order in the

Λ_h/m_B expansion with [22]

$$K_1^{\|(a)}(q^2) = -\frac{\lambda_u}{\lambda_t}\left(\frac{\overline{C}_1}{3} + \overline{C}_2\right)\delta_{qu} + \left(\overline{C}_4 + \frac{\overline{C}_3}{3}\right). \tag{2.13}$$

Here the Wilson coefficients $\overline{C}_i$ refer to the basis defined in [22] (at leading order it coincides with the basis used in [4]). The functions $K_{1,2}^{\perp}$ in (2.11) only appear at sub-leading order of the Λ_h/m_B expansion. Nevertheless, numerically the leading order effect in $b_q^{\|}(q^2)$ was found to be rather small [22], whereas the formally sub-leading effects entering $b_q^{\perp}(0)$ were found to be sizeable [26]. In particular, the contributions from the scalar penguin operators O_5 and O_6 play an important role, despite the fact that the corresponding Wilson coefficients are rather small. In the following we will extend the calculation of Kagan/Neubert in [26] and study how these effects change with increasing photon virtuality q^2. Note that in this section we will only consider those corrections that are sensitive to the charge of the spectator quark, and neglect all sub-leading effects that do *not* contribute to the isospin asymmetry. Furthermore we discard corrections to $b_q^{\|}$ that are sub-leading in Λ_h/m_B. The calculation of such effects requires a better understanding of the light-cone wave functions of the B meson and is beyond the scope of this paper. However, we have convinced ourselves that the analogous effects that enhance the isospin asymmetry for transverse polarizations are absent in the longitudinal case, see discussion after (2.14) below.

2.2.1 Details of the calculation

(a) Annihilation graphs: the annihilation graphs in figure 2(a) involve hadronic matrix elements of the four-quark operators $\mathcal{O}_{1-6}$. To first approximation, these can be factorized into two matrix elements defining meson decay constants and meson-photon transition form factors, according to the two graphs in figure 2(a). The meson-photon transition form factors and their role in the annihilation contribution to B decays have been extensively discussed in the recent literature [31] and [37]-[44].

Let us start with the matrix element of the operators $\mathcal{O}_5$ and $\mathcal{O}_6$ which can be expressed in terms of the (pseudo-)scalar $K^*\gamma^*$ transition form factor and the B meson decay constant f_B. (An analogous term with the (pseudo-)scalar $B\gamma^*$ transition form factor is absent, because the remaining matrix element of a pseudo-scalar current between a vector meson and the vacuum vanishes.) Note that the annihilation graphs with $\mathcal{O}_{5,6}$ are independent of the B meson wave function at the considered level of approximation. We calculate the (pseudo-)scalar $K^*\gamma^*$ transition form factor perturbatively in terms of the K^* light-cone wave functions, which yields

$$\begin{aligned}\langle K^*(p',\varepsilon)\gamma^*(q,\mu)|\bar{s}(1+\gamma_5)q|0\rangle &= 2\,i\,e\,e_q\,f_{K^*}^{\perp}\left((q\cdot\varepsilon^*)\,p'^{\mu} - \varepsilon^{*\mu}\,(q\cdot p') - i\epsilon^{\mu\alpha\beta\gamma}\varepsilon^*_\alpha q_\beta p'_\gamma\right)\times \\ &\times\int_0^1 du\,\frac{\phi_\perp(u)}{\bar{u}\,m_B^2 + u\,q^2} + (e_q\to e_s,\ u\to\bar{u})\,. \end{aligned}\tag{2.14}$$

Note that only transverse polarizations do contribute here. (This remains true if one includes sub-leading contributions from twist-3 wave functions. Therefore the part of the isospin asymmetry which is related to longitudinal K^* polarizations receives no enhanced contributions from the Wilson coefficients C_5 and C_6.) Comparing with (2.11) we obtain the following contribution to the function $K_1^\perp$

$$K_1^{\perp(a)}(q^2) = -\left(\overline{C}_6 + \frac{\overline{C}_5}{N_C}\right) F_\perp\left(\frac{q^2}{m_B^2}\right), \qquad F_\perp(\hat{s}) = \frac{1}{3}\int_0^1 du\,\frac{\phi_\perp(u)}{\bar{u}+u\hat{s}}\,, \tag{2.15}$$

which coincides with the result in [26] for $q^2 \to 0$. (When comparing with [26] one has to use $\lambda_c \approx -\lambda_t$. Note that in (2.15) and in the following we keep terms with $\hat{s}$ in the denominator of convolution integrals with light meson distribution amplitudes, although these terms are formally suppressed by $\hat{s}\ln\hat{s}$ for small values of $\hat{s} = q^2/m_B^2$.) The functions $K_2^\perp$ and $K_1^\parallel$ do not receive contributions from C_5 and C_6 to the considered order in the Λ_h/m_B expansion.

The $(V-A)\otimes(V-A)$ currents from the operators $\mathcal{O}_1$–$\mathcal{O}_4$ give the leading contribution to the function $K_1^\parallel$ quoted in eq. (2.13). For transverse K^* polarizations one obtains the contribution

$$K_2^{\perp(a)}(q^2) = -\frac{\lambda_u}{\lambda_t}\left(\frac{\overline{C}_1}{3} + \overline{C}_2\right)\delta_{qu} + \left(\overline{C}_4 + \frac{\overline{C}_3}{3}\right), \tag{2.16}$$

which again coincides with the term quoted in [26]. Note that the separate terms in the decomposition of $\langle K^*\gamma^*|(V-A)\otimes(V-A)|B\rangle$ are invariant under QED gauge transformations only up to a "contact term" which does not depend on the meson's wave functions [38, 39, 41]. In the sum, of course, these terms cancel.

We finally mention that there is another mechanism to induce the $B \to K^*\gamma^*$ decay where the photon does not couple in a point-like way, but enters through its hadronic component [45, 46]. This effect gives rise to $1/m_b$ corrections to the meson-photon transition form factors, which are not taken into account here. We will also neglect NLO order terms arising from α_s corrections to the annihilation topologies. In the numerical analysis below, we will use leading-order values for the Wilson coefficients C_1 and C_2, and next-to-leading order values for $C_3 - C_6$, and include the uncertainty with respect to variations of the factorization scale μ into our error estimate.

(*b*) Hard spectator interactions involving O_8: calculating the decay amplitude from the diagrams in figure 2(*b*), with the insertion of the gluonic penguin operator $\mathcal{O}_8$, we obtain

$$\begin{aligned} K_1^{\perp(b)}(q^2) &= C_8^{\text{eff}}\,\frac{m_b}{m_B}\,\frac{C_F}{N_c}\,\frac{\alpha_s}{4\pi}\,X_\perp\left(\frac{q^2}{m_B^2}\right), \\ X_\perp(\hat{s}) &= F_\perp(\hat{s}) + \frac{1}{3}\int_0^1 du\,\frac{1}{(\bar{u}+u\hat{s})^2}\,\phi_\perp(u)\,, \qquad (2.17) \\ K_2^{\perp(b)}(q^2) &= \mathcal{O}\left(\frac{\Lambda_h}{m_B}\right), \qquad (2.18) \end{aligned}$$

which again coincides with [26] in the limit $q^2 \to 0$. For the longitudinal part we have

$$K_1^{\parallel(b)}(q^2) = -C_8^{\text{eff}}\,\frac{m_b}{m_B}\,\frac{C_F}{N_c}\,\frac{\alpha_s}{4\pi}\,F_\parallel\left(\frac{q^2}{m_B^2}\right), \qquad F_\parallel(\hat{s}) = 2\int_0^1 du\,\frac{\phi_\parallel(u)}{\bar{u}+u\hat{s}}\,, \tag{2.19}$$

which has already been calculated in [22]. Note that the function $X_\perp^{(1)}(\hat{s})$ contains a logarithmic divergence for $\hat{s} \to 0$ which indicates the breakdown of QCD factorization for matrix elements with $\mathcal{O}_8$ at sub-leading order in the Λ_h/m_B expansion. We follow the authors of [26] and parameterize the divergent part of the integral by replacing

$$\int_0^1 du \to (1+\rho\, e^{i\phi}) \int_0^{1-\Lambda_h/m_B} du \tag{2.20}$$

and taking $\Lambda_h \simeq 0.5\,\mathrm{GeV}$ as a regulating cut-off, together with $0 \le \rho \le 1$ and ϕ, giving a conservative estimate of the theoretical uncertainty related to the absolute value and phase of this non-factorizing soft contribution.

(c) **Hard spectator interactions involving $\mathcal{O}_{1-6}$:** the calculation of the diagrams in figure 2(c) splits into two parts: First one calculates the internal quark loop which results in an effective $b \to sg^*$ vertex with vector- and tensor-like couplings and form factors depending on the internal quark kinematics and the quark masses in the loop. The tensor contribution is local, and can be, as usual, absorbed into $C_8^{\rm eff}$. The vector form factor is expressed in terms of the function ($z \equiv 4m_q^2/s$)

$$h(s,m_q) = -\frac{4}{9}\left(\ln\frac{m_q^2}{\mu^2} - \frac{2}{3} - z\right) - \frac{4}{9}\,(2+z)\,\sqrt{|z-1|}\,\begin{cases} \arctan\dfrac{1}{\sqrt{z-1}} & z>1 \\ \ln\dfrac{1+\sqrt{1-z}}{\sqrt{z}} - \dfrac{i\pi}{2} & z\le 1 \end{cases} \tag{2.21}$$

and reads

$$\begin{aligned} F_V(s) \;=\; \frac{3}{4}\Big\{ & h(s,m_c)\,(\overline{C}_2+\overline{C}_4+\overline{C}_6) + h(x,m_b)\,(\overline{C}_3+\overline{C}_4+\overline{C}_6) + \\ & + h(s,0)\,(\overline{C}_3+3\overline{C}_4+3\overline{C}_6) - \frac{8}{27}\,(\overline{C}_3-\overline{C}_5-15\overline{C}_6)\Big\}. \end{aligned} \tag{2.22}$$

The rest of the calculation is completely analogous as for the diagrams with $\mathcal{O}_8$ in figure 2(b). This yields the following contributions to the isospin asymmetry,

$$K_1^{\perp(c)}(q^2) = \frac{C_F}{N_c}\,\frac{\alpha_s}{4\pi}\,\frac{2}{3}\int_0^1 du\,\frac{\phi_\perp(u)}{\bar{u}+u\,\hat{q}^2}\,F_V(\bar{u}\,m_B^2+u\,q^2)\,; \tag{2.23}$$

$$K_2^{\perp(c)}(q^2) = -\frac{C_F}{N_c}\,\frac{\alpha_s}{4\pi}\,\frac{1}{2}\int_0^1 du\,\left(g_\perp^{(v)}(u) - \frac{g_\perp'^{\,(a)}(u)}{4}\right)\,F_V(\bar{u}\,m_B^2+u\,q^2); \tag{2.24}$$

for the transverse part, and

$$K_1^{\|(c)}(q^2) = -\frac{C_F}{N_c}\,\frac{\alpha_s}{4\pi}\,2\int_0^1 du\,\phi_\|(u)\,F_V(\bar{u}\,m_B^2+u\,q^2); \tag{2.25}$$

for the longitudinal one. In the limit $q^2 \to 0$ we reproduce the result in [26] (neglecting the small Wilson coefficients C_{3-6} in $F_V(s)$).

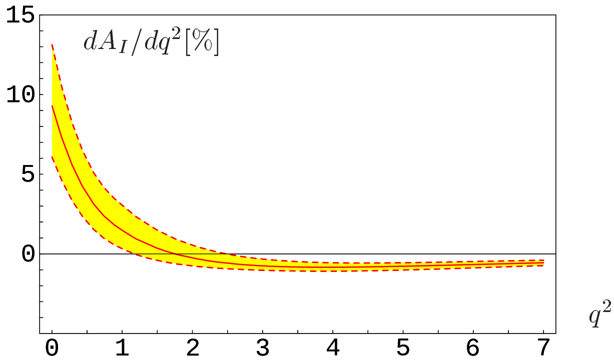

Figure 3: The SM isospin asymmetry dA_I/dq^2 for the decay $B \to K^*\ell^+\ell^-$ as a function of q^2. The error band includes the variation of hadronic and SM input parameters (taken as in [22]) and of the renormalization scale μ between $m_b/2$ and $2m_b$.

2.2.2 Numerical estimate

We now turn to the numerical analysis of the isospin asymmetry in $B \to K^*\ell^+\ell^-$, combining the above results as

$$\begin{aligned} K_1^\perp(q^2) &= K_1^{\perp(a)}(q^2) + K_1^{\perp(b)}(q^2) + K_1^{\perp(c)}(q^2)\,, \\ K_2^\perp(q^2) &= K_2^{\perp(a)}(q^2) + K_2^{\perp(c)}(q^2)\,, \\ K_1^{||}(q^2) &= K_1^{||(a)}(q^2) + K_1^{||(b)}(q^2) + K_1^{||(c)}(q^2)\,. \end{aligned} \tag{2.26}$$

The central values and uncertainties of all hadronic and SM input parameters are taken as in ref. [22]. The renormalization scale in the annihilation and hard-scattering processes will be taken as $\mu' = \sqrt{\mu\Lambda_h}$, where μ will be varied as $m_b/2 \le \mu \le 2m_b$.

The result for the isospin asymmetry is plotted as a function of q^2 in figure 3. At $q^2 = 0$ we basically recover the result of Kagan/Neubert [26] with

$$\mathrm{Re}\left[b_d^\perp(0) - b_u^\perp(0)\right] = +\,9.3\,{}^{+3.8}_{-3.2}\,\%\,, \tag{2.27}$$

where the uncertainty due to the theoretical input parameters (taken as in [22]) is dominated by the soft form factor $\xi_\perp$ ($\sim {}^{+2.3}_{-1.5}\%$), the B meson decay constant ($\sim 1.5\%$), the moment of the B meson light-cone wave function λ_B ($\sim 1.1\%$), and the IR-sensitive integral $X_\perp^{(1)}$ ($\sim 1\%$). We thus confirm the conclusion in [26] that QCD factorization correctly reproduces the sign and magnitude of the experimentally measured isospin asymmetry ($A_I[B \to K^*\gamma] = 0.11 \pm 0.07$ [47]–[49]) with SM values for the Wilson coefficients.

For increasing values of q^2 the isospin-asymmetry decreases, and its central value becomes slightly negative above $q^2 = 2\,\mathrm{GeV}^2$ and stays basically at a constant value of about -1%. Since the uncertainty related to the hadronic input parameters is reduced as well, this means that the measurement of a significant deviation from zero of the isospin asymmetry in the range $2\,\mathrm{GeV}^2 < q^2 < 7\,\mathrm{GeV}^2$ may still indicate new physics (although one would need to have a handle on even higher order effects, before drawing any definite conclusions). Note that in the SM the isospin asymmetry is sensitive to C_5 and C_6 at small q^2 but to C_3 and C_4 at larger q^2. Thus, in principle, the two momentum regions provide complementary tests of the four-quark penguin-operators (see also figure 6 below).

3. Decay asymmetries in MSSM with MFV

In this section we present an example to illustrate how new physics may affect the FB and isospin asymmetry in $B \to K^* \ell^+ \ell^-$ decay. We consider the minimal supersymmetric standard model (MSSM) with minimal flavor violation (MFV) [10, 50, 51]. In this model, all flavour transitions are governed by the CKM matrix. In addition, in order to increase the predictivity of the model, we impose the following restrictions.

- All supersymmetric particles are taken as heavy (about 1 TeV), except for charginos, sneutrinos, the light (mostly right-handed) stop, and charged Higgs fields. Heavy particles (squarks, gluinos etc.) are integrated out to obtain a "low-energy" effective theory in terms of light SUSY particles and SM fields [10].
- The weak effective hamiltonian (2.1) includes only SM operators.
- We assume a flavor diagonal mass matrix in the down-squark sector [9]. Therefore neutralino contributions to the Wilson coefficients will be absent.

We carefully consider the two regimes of small and large $\tan\beta$, taking into account the modifications and subtleties associated with large values of $\tan\beta$ [11, 52] ($\tan\beta$ is the ratio of vacuum expectation values of the two Higgs doublets).

The analysis is organized in three steps: first, we perform a scan of model parameters within specific ranges which will be indicated below. For each set of parameter values we will calculate the Wilson coefficients contributing to the $B \to X_s\gamma$ decay at NLO following [10, 11, 52], including also electroweak corrections [3] and non-perturbative contributions [5]. Then we will use the experimental data on $B \to X_s\gamma$ to constrain the model parameter space, i.e. those parameter sets that lead to Wilson coefficients which are incompatible with the $B \to X_s\gamma$ constraint will be discarded. (Note that the allowed region of parameter space will divide into two classes, corresponding to negative or positive sign of $C_7^{\rm eff}$.) Finally, using the remaining part of parameter space we evaluate the Wilson coefficients in (2.1), in terms of which we obtain numerical results for the two decay asymmetries, $dA_{\rm FB}$ and dA_I in MFV.

3.1 Model set-up

The free parameters of the model are:

- The Higgs mixing parameter μ. We consider $|\mu| \leq 500\,\mathrm{GeV}$.
- The mass parameter M_2 in the chargino mass matrix ("weak gaugino mass"), which is varied as $M_2 = 100 - 800\,\mathrm{GeV}$. Alternatively, one could consider one of the

chargino mass eigenstates. In the parameter scan we always impose the constraint that the masses of both charginos should be heavier than 100 GeV, as required by direct searches.

- The light stop mass is taken in the range $100 \leq \tilde{m}_{t2} \leq 350\,\text{GeV}$. (Note that depending on the relative size of stop, gluino and other squark masses, one has to be careful when evaluating the Wilson coefficients. In our case, we are considering a scenario with a very light right-handed stop and charginos in comparison with gluino and other squark masses. This implies that we have to use the Wilson coefficients at the weak scale as described in [11, eq.(31)].)

- $|\theta_t|$ is the stop mixing angle, such that $\tilde{t}_1 = \cos\theta_t \tilde{t}_L + \sin\theta_t \tilde{t}_R$ and $\tilde{t}_2 = -\sin\theta_t \tilde{t}_L + \cos\theta_t \tilde{t}_R$. Constraints from electroweak precision measurements favor the light stop quark to be almost right-handed [53]. We therefore take $|\theta_t| \leq \pi/10$.

- The charged Higgs mass is considered in the range $m_H = 100 - 400\,\text{GeV}$. We have chosen a not too heavy mass for the charged Higgs to evaluate the maximal possible impact in our observables of a non-decoupled charged Higgs.

- The ratio of vacuum expectation values for the two Higgs doublets is varied as $\tan\beta = 2 - 40$. In the expressions for the Wilson coefficients the large $\tan\beta$ regime is treated as in ref. [11]. Large values of $\tan\beta$ are particularly interesting since this can easily lead to a sign flip in the value of C_7^{eff}.

The following parameters will be kept fixed in the analysis:

- According to the effective theory approximation of the MFV supersymmetric model, all heavy squark masses (including the heavy stop) are taken at a common value, $m_{\tilde{q}} = 1\,\text{TeV}$. Also the gluino mass is taken at a fixed value $m_{\tilde{g}} = 1\,\text{TeV}$. The renormalization effects from integrating out squarks and gluinos between the scale 1 TeV and the electroweak scale as discussed in [11] are taken into account.

- For the sneutrino masses we consider a conservative value, $m_{\tilde{\nu}} = 100\,\text{GeV}$ which is not in conflict with direct searches.

Scanning over discrete points in the so-defined parameters space, we calculate the Wilson coefficients for the operators in (2.1) at the electroweak scale, $\mu_W = \mathcal{O}(M_W)$. At NLO accuracy one obtains the generic form

$$\begin{aligned} C_i(\mu_W) &= C_i(\mu_W)_{\text{SM}} + \delta C_i(\mu_W) \\ &= C_i^{(0)}(\mu_W)_{SM} + \delta C_i^{(0)}(\mu_W)_H + \delta C_i^{(0)}(\mu_W)_{\text{SUSY}} + \\ &\quad + \frac{\alpha_s(\mu_W)}{4\pi}\left(C_i^{(1)}(\mu_W)_{SM} + \delta C_i^{(1)}(\mu_W)_H + \delta C_i^{(1)}(\mu_W)_{\text{SUSY}}\right), \end{aligned} \tag{3.1}$$

where we indicated explicitely the contributions from different virtual particles in box and penguin diagrams:

- SM contributions from charged gauge bosons and up-type quarks.

- Contributions from charged Higgs bosons and up-type quarks, present in SUSY, but also in generic models with two Higgs doublets.

- Specific SUSY contributions from charginos and up-type squarks, gluinos and down-type squarks, neutralinos and down-type quarks. (We will mainly focus on chargino effects. In our model gluinos are very heavy and their main effect is to induce renormalization effects in the vertices involving charginos, charged Higgs bosons, charged gauge bosons and unphysical scalars [10]. As stated above, neutralino contributions are absent due to our assumption about the mass matrix in the down-quark sector.)

The expressions for $\delta C_{7,8}^{\mathrm{eff}\,(0,1)}$ can be found in ref. [10]; SUSY contributions to C_9 and C_{10} are taken from [13, 14]. We have also included the NLO SUSY contribution to the four-quark penguin operators from ref. [9]. (Notice the different conventions used for θ_t and the sign of the μ parameter in [10] and [9].) The values of the Wilson coefficients are evolved to the low scales relevant for the exclusive B decays by standard renormalization group analysis.

3.2 Constraints from $B \to X_s\gamma$

We consider the world average of measurements of the inclusive branching ratio for $B \to X_s\gamma$ [17]

$$2.82 \times 10^{-4} \leq \mathrm{BR}(B \to X_s\gamma) \leq 3.62 \times 10^{-4} \,. \tag{3.2}$$

The precision reached in the determination of this decay mode significantly constrains new physics contributions and requires to consider very carefully higher order QCD effects in the theoretical calculations, also in the presence of new physics. For instance, at two loops the electroweak amplitude gets enhanced in the SM due to a large QCD mixing with the effective operator $b \to s\bar{c}c$. Moreover, in the MFV model we are discussing it is known that there is a compensatory effect between the charged Higgs and chargino-stop contribution that gets unbalanced once the NLO correction is included. Apart from the standard renormalization group improvement of Wilson coefficients at NLO, we finally have to include non-perturbative corrections associated with the hadronic matrix elements for $B \to X_s\gamma$. In the analysis we will strictly follow refs. [10, 11].

Scanning over the allowed parameter space (notice that in our scenario we are only exploring light stop, charginos and not too heavy charged Higgses), we find the following allowed range for the new contributions δC_7 to the Wilson coefficient C_7^{eff} at the scale μ_W (see eq. (3.1))

$$-0.02 \leq \delta C_7(\mu_W) \leq 0.12 \qquad \text{or} \qquad 0.59 \leq \delta C_7(\mu_W) \leq 1.24 \,, \tag{3.3}$$

where the first range is SM-like (negative sign for C_7^{eff}), and the second range corresponds to the flipped-sign solution. The behaviour of the two ranges as a function of the parameters $\tan\beta$, θ_t and $\tilde{m}_{t2}$ is quite different. We have found that the points in the SM-like region corresponds predominantly to values for $\tan\beta$ in the low regime (between 2 to 10). Moreover, higher values of m_h tend to prefer lower values of $\tan\beta$. There is no preferred

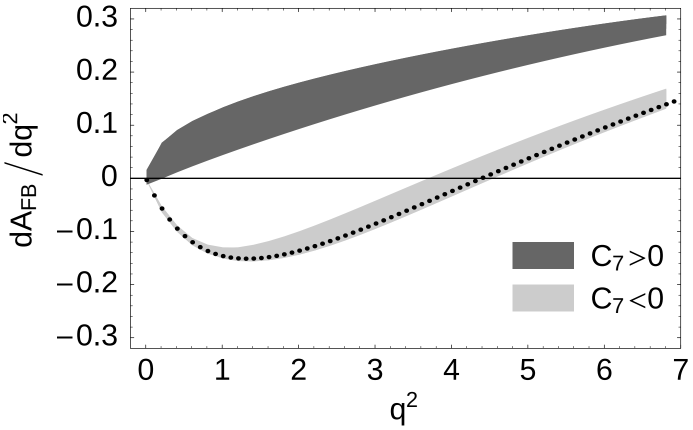

Figure 4: FB asymmetry (2.2) for the decay $B \to K^* \ell^+ \ell^-$ as a function of q^2. The light band corresponds to the region of parameter space with same (negative) sign for C_7 as in the SM; the dark band refers to solutions with flipped sign for C_7; and the dotted line is the central value of the SM prediction.

value for θ_t and also any value of $\tilde{m}_{t2}$ is found. The flipped-sign solutions for $C_7^{\rm eff}$ correspond to large values of $\tan\beta$ of order 40. In this case, θ_t preferably takes extreme values around $\pm\pi/10$. The allowed values for $\tilde{m}_{t2}$ increase with increasing m_h but are always below 300 GeV, for example, for $m_h = 100\,{\rm GeV}$ one gets $\tilde{m}_{t2} = 100 - 240\,{\rm GeV}$, and for the highest value $m_h = 400\,{\rm GeV}$ one has $\tilde{m}_{t2} = 100 - 300\,{\rm GeV}$. The corresponding range for the Wilson coefficient $C_8^{\rm eff}$ reads

$$-0.18 \leq \delta C_8(\mu_W) \leq 0.03 \qquad \text{or} \qquad 0.20 \leq \delta C_8(\mu_W) \leq 0.57\,, \tag{3.4}$$

Concerning the Wilson coefficients $C_3 - C_6$ of four-quark penguin operators our MFV model is SM-like. Also the Wilson coefficients C_9 and C_{10} only receive tiny contributions (of at most 10% from their SM value, we restricted ourselves to LO SUSY contributions here). The mechanisms that enhance the SUSY contributions to $C_7^{\rm eff}$ at large $\tan\beta$ are not working for C_9 and C_{10}. For instance, the charged Higgs contribution dominant for $C_7^{\rm eff}$ at large $\tan\beta$ is suppressed for $C_{9,10}$ [13, 14]. Therefore the modifications for the FB and isospin-asymmetry, to be discussed below, will be due to new physics contributions to $C_7^{\rm eff}$, $C_8^{\rm eff}$ mainly.

3.3 FB and isospin asymmetry in MFV

Using the parameter values in the MFV scenario allowed from the constraints discussed in the previous subsection, we calculate the FB and isospin-asymmetry as explained in section 2. The result for the FB asymmetry is shown in figure 4. The two distinct bands refer to the two possible signs of $C_7^{\rm eff}$ in (3.3). Comparing with the SM case in figure 1, we observe that the variation of MFV parameters for the case of SM-like sign of $C_7^{\rm eff}$ (light band) is indistinguishable from the hadronic uncertainties. An exception is the FB asymmetry zero where the hadronic uncertainty is reduced due to the form factor relations discussed in [22]–[24]. Here the MFV effects are competitive with the hadronic uncertainty, and show a tendency towards lower values of the FB asymmetry zero. This implies that

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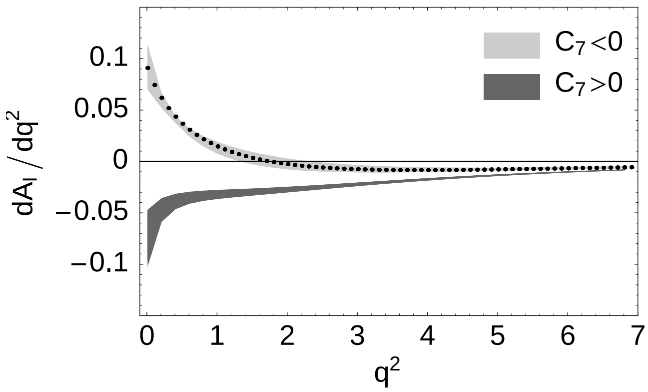

Figure 5: Isospin asymmetry (2.3) for the decay $B \to K^* \ell^+ \ell^-$ as a function of q^2. Notations as in figure 4.

if the experimental measurement of the FB asymmetry zero found an even *larger* q^2 value than predicted in the SM it would not easily be accommodated by our MFV scenario, and eventually would lead to the exclusion of the region of parameter space described in section 3.1. As is well known, for the flipped sign solution of $C_7^{\rm eff}$ (dark band) the characteristic features of the FB asymmetry change dramatically: dA_{FB}/dq^2 is positive for small *and* large values of q^2, and there is no asymmetry zero in the spectrum.

The MFV prediction for the isospin asymmetry, shown in figure 5, also splits into two bands, according to the sign of $C_7^{\rm eff}$. Because the isospin asymmetry is dominated by the penguin coefficients C_{3-6} which are only slightly modified in our MFV scenario, both bands are very thin. Within the MFV the isospin asymmetry at small values of q^2 thus provides a possibility to exclude the flipped sign solution for $C_7^{\rm eff}$, independently from the analysis of the FB asymmetry (as already mentioned, at present, experimental data favor the same-sign solution). On the other hand, a measurement of a large isospin asymmetry at moderate values of q^2 would point to physics beyond both, the SM and MFV.

4. Beyond MFV

As we have seen in the previous section, in MFV the main effect on the FB and isospin asymmetry is due to a possible sign flip in the value of $C_7^{\rm eff}$. Taking the central value of the experimental measurement of the isospin asymmetry in $B \to K^*\gamma$ at face value, one is tempted to consider this scenario as already excluded. (A reduction of theoretical and experimental uncertainties for this observable, combined with a measurement of the FB asymmetry would allow for an ultimate conclusion.) The Wilson coefficients for semi-leptonic and 4-quark penguin operators are only slightly affected in our MFV scenario, and therefore for SM-like sign of $C_7^{\rm eff}$ MFV is indistinguishable from the SM within the theoretical uncertainties.

On the other hand, in a generic new physics model, one might expect sizeable corrections to other Wilson coefficients as well, which may show up in the $B \to K^* \ell^+ \ell^-$ decay asymmetries. In the following we perform a brief discussion on the possible size of such effects.

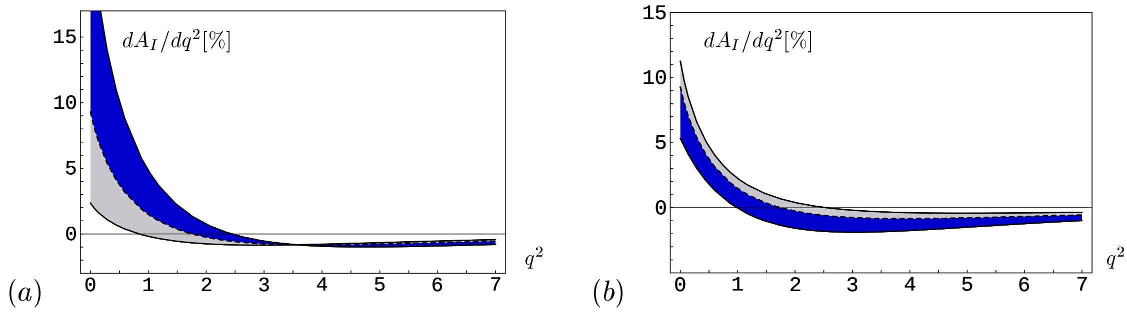

Figure 6: The isospin asymmetry dA_I/dq^2 for the decay $B \to K^*\ell^+\ell^-$ as a function of q^2. (a) The combination $a_6^{(0)} = (\bar{C}_6 + \bar{C}_5/3)$ in the function $K_1^{\perp(a)}$ is varied within a factor of two around its SM value (dark band = larger values, light band = smaller values). (b) The same for the combination $a_4^{(0)} = (\bar{C}_4 + \bar{C}_3/3)$ in the functions $K_1^{\|(a)}$ and $K_2^{\perp(a)}$.

- The penguin operators $\mathcal{O}_{3-6}$ give the main effect to the isospin asymmetry in $B \to K^*\gamma$ and $B \to K^*\ell^+\ell^-$. More precisely, dA_I/dq^2 is sensitive to the combinations $a_6^{(0)} = (\bar{C}_6+\bar{C}_5/3)$ and $a_4^{(0)} = (\bar{C}_4+\bar{C}_3/3)$ entering the functions $K_{1,2}^{\perp(a)}$ and $K_1^{\|(a)}$ that describe the dominant annihilation contribution to the isospin asymmetry at small and large q^2, respectively. In principle, the size of the corresponding Wilson coefficients is already constrained from many sources, for instance non-leptonic B decays, life-time differences of B mesons, kaon decays etc. Since the sensitivity to individual Wilson coefficients in these decays is not very high, and considering the possibility of some conspiracy between different new physics effects one may, however, still allow for sizeable deviations of C_{3-6} from their SM values. Figure 6 shows the effect on the isospin asymmetry if we allow for a variation of $a_6^{(0)}$ and $a_4^{(0)}$ within a factor of two. In this case, as already mentioned above, modifications of $a_6^{(0)}$ (respectively $a_4^{(0)}$) are competitive with the theoretical uncertainties at small (respectively large) values of q^2 (at large q^2 values, however, the isospin asymmetry itself is rather small, and probably the necessary accuracy to see modifications of $a_4^{(0)}$ will be hard to achieve.)

- The semi-leptonic operators $\mathcal{O}_{9,10}$ are relevant for $B \to K^*\ell^+\ell^-$ at not too small values of q^2. As already stated above the FB asymmetry provides a rather clean test of the Wilson coefficients $C_{9,10}$ (compared to this the isospin asymmetry is less sensitive to C_9 and C_{10}, unless the experimental measurement leads to an unexpectedly large asymmetry for $2\,\mathrm{GeV}^2 < q^2 < 7\,\mathrm{GeV}^2$.) In particular, when NLO QCD effects are included the Wilson coefficient C_9 can be measured with 10% accuracy from the FB asymmetry zero [22] (provided the relative signs of $C_{7,9,10}$ are SM-like). If the relative sign between $C_7^{\rm eff}$ and C_9 is flipped there will be no FB asymmetry zero at all; in case of a non-SM-like sign of C_{10} the overall sign of the FB asymmetry will change. We refrain from giving quantitative results here, but rather refer to [15, 24, 54] for extensive phenomenological analyses of these issues.

- In this work we do not entertain the possibility of introducing new operators in the effective hamiltonian (2.1). An example of contributions of non-standard operators to the FB asymmetry in $B \to K^* \ell^+ \ell^-$ can be found in [55]. A discussion of non-standard penguin operators, which would also be relevant for the isospin asymmetry, has been given in [56].

5. Summary

We discussed two dedicated observables in $B \to K^* \ell^+ \ell^-$ decays, the forward-backward and isospin-asymmetry, which can be used to test the standard model against new physics. We concentrated on the large recoil region where the decay amplitudes can be calculated using the QCD factorization approach. (For the FB asymmetry we used the results from [22]; for the isospin asymmetry we generalized the $B \to K^* \gamma$ analysis performed in [26].) The two asymmetries turn out to give complementary tests of the Wilson coefficients in the weak effective hamiltonian: For a given value (and sign) of $C_7^{\rm eff}$, the FB asymmetry is sensitive to the semi-leptonic coefficients C_9 and C_{10}. On the other hand the isospin asymmetry is dominated by the effects of four-quark coefficients $C_{5,6}$ and $C_{3,4}$ at small or medium values of q^2.

To illustrate possible new physics effects we considered an explicit model, namely minimal supersymmetry with minimal flavor violation. In this scenario the main deviation from the SM is the possibility to generate another sign for the magnetic penguin coefficients $C_7^{\rm eff}$, $C_8^{\rm eff}$ (while the absolute value for $C_7^{\rm eff}$ is constrained by the $B \to X_s \gamma$ measurement). The corresponding two regions of allowed parameter sets translate into two well separated bands of predictions for the FB and isospin-asymmetry. On the other hand, the change in the remaining Wilson coefficients is rather small in the MFV scenario and indistinguishable from hadronic uncertainties of the theoretical calculation.

Considering an alternative point of view and assuming the SM to be a valid approximation in rare B decays, the simultaneous study of the isospin asymmetry and the forward-backward asymmetry in the large-recoil region ($q^2 < 4m_c^2$) provides a perfect testing ground for the QCD factorization approach.

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