Looking for the Unexpected: 
Direct CP Violation in $B \to X_s \gamma$ Decays

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Abstract:
The observation of a sizable direct CP asymmetry in the inclusive decays $B \to X_s \gamma$ would be a clean signal of New Physics. In the Standard Model, this asymmetry is below 1% in magnitude. In extensions of the Standard Model with new CP-violating couplings, large asymmetries are possible without conflicting with the experimental value of the $B \to X_s \gamma$ branching ratio. In particular, large asymmetries arise naturally in models with enhanced chromo-magnetic dipole transitions. Some generic examples of such models are explored and their implications for the semileptonic branching ratio and charm yield in $B$ decays discussed.

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DIRECT CP VIOLATION IN $B \rightarrow X_s\gamma$ DECAYS

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The observation of a sizable direct CP asymmetry in the inclusive decays $B \rightarrow X_s\gamma$ would be a clean signal of New Physics. In the Standard Model, this asymmetry is below 1% in magnitude. In extensions of the Standard Model with new CP-violating couplings, large asymmetries are possible without conflicting with the experimental value of the $B \rightarrow X_s\gamma$ branching ratio. In particular, large asymmetries arise naturally in models with enhanced chromo-magnetic dipole transitions. Some generic examples of such models are explored and their implications for the semileptonic branching ratio and charm yield in $B$ decays discussed.

1 Introduction

Studies of rare decays of $B$ mesons have the potential to uncover the origin of CP violation and provide hints to physics beyond the Standard Model of strong and electroweak interactions. The measurements of several CP asymmetries will make it possible to test whether the CKM paradigm is correct, or whether additional sources of CP violation are required. In order to achieve this goal, it is necessary that the theoretical calculations of CP-violating observables in terms of Standard Model parameters are, to a large extent, free of hadronic uncertainties. This can be achieved, e.g., by measuring time-dependent asymmetries in the decays of neutral $B$ mesons into particular CP eigenstates. In many other cases, however, the theoretical predictions for direct CP violation in exclusive $B$ decays are obscured by large strong-interaction effects.

Inclusive decay rates of $B$ mesons, on the other hand, can be reliably calculated in QCD using the operator product expansion. Up to small bound-state corrections these rates agree with the parton model predictions for the underlying decays of the $b$ quark. The disadvantage that the sum over many final states partially dilutes the CP asymmetries in inclusive decays is compensated by the fact that, because of the short-distance nature of these processes, the strong phases are calculable using quark–hadron duality. In this talk, I report on a study of direct CP violation in the rare radiative decays $B \rightarrow X_s\gamma$, both in the Standard Model and beyond. These decays have already been observed experimentally, and copious data samples will be collected at the $B$ factories. The theoretical analysis relies only on the weak assumption of global quark–hadron duality, and the leading nonperturbative corrections are well understood.
We perform a model-independent analysis of CP-violating effects in $B \to X_s \gamma$ decays in terms of the effective Wilson coefficients $C_7 \equiv C_7^{\text{eff}}(m_b)$ and $C_8 \equiv C_8^{\text{eff}}(m_b)$ multiplying the (chromo-) magnetic dipole operators $O_7 = e m_b \bar{s}_L \gamma_\mu F^{\mu \nu} b_R$ and $O_8 = g_s m_b \bar{s}_L \sigma_{\mu \nu} G^{\mu \nu} b_R$ in the effective weak Hamiltonian. We allow for generic New Physics contributions to these coefficients. Several extensions of the Standard Model in which such contributions arise have been explored, e.g., in Refs. 7–10. We find that in the Standard Model the direct CP asymmetry in $B \to X_s \gamma$ decays is very small (below 1% in magnitude) because of a combination of CKM and GIM suppression, both of which can be lifted in New Physics scenarios with additional contributions to the dipole operators containing new weak phases. We thus propose a measurement of the inclusive CP asymmetry in radiative $B$ decays as a clean and sensitive probe of New Physics. Studies of direct CP violation in the inclusive decays $B \to X_s \gamma$ have been performed previously by several authors, both in the Standard Model and in certain extensions of it. In all cases rather small asymmetries were obtained. We generalize and extend these analyses in various ways. Besides including some contributions neglected in previous works, we investigate a class of New Physics models with enhanced chromo-magnetic dipole contributions, in which large CP asymmetries of order 10–50% are possible and even natural. We also employ a full next-to-leading order analysis of the CP-averaged $B \to X_s \gamma$ branching ratio in order to derive constraints on the parameter space of the New Physics models considered.

## 2 Direct CP violation in radiative $B$ decays

The starting point in the calculation of the inclusive $B \to X_s \gamma$ decay rate is provided by the effective weak Hamiltonian renormalized at the scale $\mu = m_b$. Direct CP violation in these decays may arise from the interference of non-trivial weak phases, contained in CKM parameters or in possible New Physics contributions to the Wilson coefficient functions, with strong phases provided by the imaginary parts of the matrix elements of the operators in the effective Hamiltonian. These imaginary parts first arise at $O(\alpha_s)$ from loop diagrams containing charm quarks, light quarks or gluons. Using the formulae of Greub et al. for these contributions, we calculate at next-to-leading order the difference $\Delta \Gamma = \Gamma(B \to X_s \gamma) - \Gamma(B \to \bar{X_s} \gamma)$ of the CP-conjugate, inclusive decay rates. The contributions to $\Delta \Gamma$ from virtual corrections arise from interference of the one-loop diagrams with insertions of the operators $O_2$ and $O_8$ with the tree-level diagram containing $O_7$. Here $O_2 = \bar{s}_L \gamma_\mu q_L \bar{q}_L \gamma^{\mu} b_L$ with $q = c, u$ are the usual current–current operators in the effective Hamiltonian. There are also contributions to $\Delta \Gamma$ from gluon bremsstrahlung diagrams with a
charm-quark loop. They can interfere with the tree-level diagrams for $b \to s\gamma g$ containing an insertion of $O_7$ or $O_8$. Contrary to the virtual corrections, for which in the parton model the photon energy is fixed to its maximum value, the gluon bremsstrahlung diagrams lead to a non-trivial photon spectrum, and so the results depend on the experimental lower cutoff on the photon energy.

We define a quantity $\delta$ by the requirement that $E_{\gamma} > (1 - \delta)E_{\gamma}^{\text{max}}$. Combining the two contributions and dividing the result by the leading-order expression for twice the CP-averaged inclusive decay rate, we find for the CP asymmetry

$$A_{\text{CP}}^{b \to s\gamma}(\delta) = \frac{\Gamma(B \to X_s\gamma) - \Gamma(B \to X_{\bar{s}}\gamma)}{\Gamma(B \to X_s\gamma) + \Gamma(B \to X_{\bar{s}}\gamma)} \bigg|_{E_{\gamma} > (1 - \delta)E_{\gamma}^{\text{max}}}$$

$$= \frac{\alpha_s(m_b)}{|C_7|^2} \left\{ \frac{40}{81} \text{Im}[C_2C_7^*] - \frac{8z}{9} \left[ v(z) + b(z, \delta) \right] \text{Im}[(1 + \epsilon_s)C_2C_7^*] \right. - \left. \frac{4}{9} \text{Im}[C_8C_7^*] + \frac{8z}{27} b(z, \delta) \text{Im}[(1 + \epsilon_s)C_2C_8^*] \right\},$$

where $z = (m_c/m_b)^2$, and the explicit expressions for the functions $g(z)$ and $b(z, \delta)$ can be found in Ref. 6. The quantity $\epsilon_s$ is a ratio of CKM matrix elements given by

$$\epsilon_s = \frac{V_{us}^* V_{ub}}{V_{ts}^* V_{tb}} \approx \lambda^2(i\eta - \rho) = O(10^{-2}),$$

where $\lambda = \sin\theta_C \approx 0.22$ and $\rho, \eta = O(1)$ are the Wolfenstein parameters. An estimate of the $C_2$–$C_7$ interference term in $\epsilon_s$ was obtained previously by Soares 11, who neglects the contribution of $b(z, \delta)$ and uses an approximation for the function $v(z)$. The relevance of the $C_8$–$C_7$ interference term for two-Higgs-doublet models, and for left–right symmetric extensions of the Standard Model, has been explored in Refs. 12,13.

In the Standard Model, the Wilson coefficients take the real values $C_2 \approx 1.11$, $C_7 \approx -0.31$ and $C_8 \approx -0.15$. The imaginary part of the small quantity $\epsilon_s$ is thus the only source of CP violation. Note that all terms involving this quantity are GIM suppressed by a power of the small ratio $z = (m_c/m_b)^2$, reflecting the fact that there is no non-trivial weak phase difference in the limit where $m_c = m_u = 0$. Hence, the Standard Model prediction for the CP asymmetry is suppressed by three small factors: $\alpha_s(m_b)$ arising from the strong phases, $\sin^2\theta_C$ reflecting the CKM suppression, and $(m_c/m_b)^2$ resulting from the GIM suppression. The numerical result for the asymmetry depends on the values of the strong coupling constant and the ratio of the heavy-quark
Table 1: Values of the coefficients $a_{ij}$ in %, without (left) and with (right) Fermi motion effects included

| $\delta$ | $E_{\gamma}^{\text{min}}$ [GeV] | $a_{27}$ | $a_{87}$ | $a_{28}$ | $a_{27}$ | $a_{87}$ | $a_{28}$ |
|---------|---------------------------------|----------|----------|----------|----------|----------|----------|
|         | (parton model)                  |          |          |          | (with Fermi motion) |          |          |          |
| 1.00    | 0.00                            | 1.06     | -9.52   | 0.16     | 1.06     | -9.52   | 0.16     |
| 0.30    | 1.85                            | 1.17     | -9.52   | 0.12     | 1.23     | -9.52   | 0.10     |
| 0.15    | 2.24                            | 1.31     | -9.52   | 0.07     | 1.40     | -9.52   | 0.04     |

pole masses, for which we take $\alpha_s(m_b) \approx 0.214$ (corresponding to $\alpha_s(m_Z) = 0.118$ and two-loop evolution down to the scale $m_b = 4.8 \text{ GeV}$) and $\sqrt{z} = m_c/m_b = 0.29$. This yields $A^{b \to s\gamma}_{\text{CP,SM}} \approx (1.5–1.6)\% \eta$ depending on the value of $\delta$. With $\eta \approx 0.2–0.4$ as suggested by phenomenological analyses, we find a tiny asymmetry of about 0.5%, in agreement with the estimate obtained in Ref. 11. Expression (3) applies also to the decays $B \to X_d\gamma$, the only difference being that in this case the quantity $\epsilon_s$ must be replaced with the corresponding quantity $\epsilon_d = (V_{ud} V_{ub})/(V_{td} V_{tb}) \approx (\rho - i\eta)/(1 - \rho + i\eta) = O(1)$. Therefore, in the Standard Model the CP asymmetry in $B \to X_d\gamma$ decays is larger by a factor of about 20 than that in $B \to X_s\gamma$ decays. However, experimentally it is difficult to distinguish between $B \to X_d\gamma$ and $B \to X_d\gamma$ decays. If only their sum is measured, the CP asymmetry vanishes by CKM unitarity.

From (3) it is apparent that two of the suppression factors operative in the Standard Model, $z$ and $\Lambda^2$, can be avoided in models where the effective Wilson coefficients $C_7$ and $C_8$ receive additional contributions involving non-trivial weak phases. Much larger CP asymmetries then become possible. In order to investigate such models, we may to good approximation neglect the small quantity $\epsilon_s$ and write

$$A^{b \to s\gamma}_{\text{CP}}(\delta) = a_{27}(\delta) \text{Im} \left[ \frac{C_2}{C_7} \right] + a_{87} \text{Im} \left[ \frac{C_8}{C_7} \right] + a_{28}(\delta) \text{Im} \left[ \frac{C_2 C_8^*}{|C_7|^2} \right]. \quad (3)$$

The values of the coefficients $a_{ij}$ are shown in the left portion of Table 1 for three choices of the cutoff on the photon energy: $\delta = 1$ corresponding to the (unrealistic) case of a fully inclusive measurement, $\delta = 0.3$ corresponding to a restriction to the part of the spectrum above 1.85 GeV, and $\delta = 0.15$ corresponding to a cutoff that removes almost all of the background from $B$ decays into charmed hadrons. In practice, a restriction to the high-energy part of the photon spectrum is required for experimental reasons. Note, however, that the result for the CP asymmetry is not very sensitive to the choice of the photon-energy cutoff. Whereas the third term in (3) is generally very small,
the first two terms can give rise to sizable effects. Assume, e.g., that there is a New Physics contribution to $C_7$ of similar magnitude as the Standard Model contribution (so as not to spoil the prediction for the $B \rightarrow X_s \gamma$ branching ratio) but with a non-trivial weak phase. Then the first term in (3) may give a contribution of up to about 5% in magnitude. Similarly, if there are New Physics contributions to $C_7$ and $C_8$ such that the ratio $C_8/C_7$ has a non-trivial weak phase, the second term may give a contribution of up to about $10\% \times |C_8/C_7|$. In models with a strong enhancement of $|C_8|$ with respect to its Standard Model value, there is thus the possibility of generating a large direct CP asymmetry in $B \rightarrow X_s \gamma$ decays.

In our discussion so far we have neglected nonperturbative power corrections to the inclusive decay rates. Their impact on the rate ratio defining the CP asymmetry is very small, since most of the corrections cancel between the numerator and the denominator. Potentially the most important bound-state effect is the Fermi motion of the $b$ quark inside the $B$ meson, which determines the shape of the photon energy spectrum in the endpoint region. This effect is included in the heavy-quark expansion by resumming an infinite set of leading-twist corrections into a “shape function”, which governs the momentum distribution of the heavy quark inside the meson $^{14,15}$. The physical decay distributions are obtained from a convolution of parton model spectra with this function. In the process, phase-space boundaries defined by parton kinematics are transformed into the proper physical boundaries defined by hadron kinematics. Details of the implementation of this effect can be found in Refs. 6,18, where a simple ansatz for the shape function is employed. The right portion of Table 1 shows the values of the coefficients $a_{ij}(\delta)$ corrected for Fermi motion. The largest coefficient, $a_{87}$, is not affected, and the impact on the other two coefficients is rather mild. As a consequence, the predictions for the CP asymmetry are quite insensitive to bound-state effects, even if a restriction on the high-energy part of the photon spectrum is imposed.

In the next section we explore the structure of New Physics models with a potentially large inclusive CP asymmetry. A non-trivial constraint on such models is that they must yield an acceptable result for the total, CP-averaged $B \rightarrow X_s \gamma$ branching ratio, which has been measured experimentally. Taking a weighted average of the results reported by the CLEO and ALEPH Collaborations $^{19,20}$ gives $B(B \rightarrow X_s \gamma) = (2.5 \pm 0.6) \times 10^{-4}$. The complete theoretical prediction for the $B \rightarrow X_s \gamma$ branching ratio at next-to-leading order has been presented for the first time by Chetyrkin et al. $^{21}$ and subsequently has been discussed by several authors $^{22-24}$. It depends on the Wilson coefficients $C_2$, $C_7$ and $C_8$ through the combinations $\text{Re}[C_6 C_7^*]$. Recently, we have extended these analyses in several aspects, including a discussion of Fermi motion effects

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and a conservative analysis of truncation errors. In contrast to the case of the CP asymmetry, Fermi motion effects are very important when comparing experimental data for the $B \to X_s \gamma$ branching ratio with theoretical predictions. With our choice of parameters, we obtain for the total branching ratio in the Standard Model $B(\to X_s \gamma) = (3.3 \pm 0.3) \times 10^{-4}$, which is consistent with the experimental findings.

3 CP asymmetry beyond the Standard Model

In order to explore the implications of various New Physics scenarios for the CP asymmetry and branching ratio in $B \to X_s \gamma$ decays it is useful to express the Wilson coefficients $C_7 = C_{7\text{eff}}(m_b)$ and $C_8 = C_{8\text{eff}}(m_b)$, which are defined at the scale $m_b$, in terms of their values at the high scale $m_W$, using the renormalization group. This yields

\[ C_7 = \eta^{\frac{16}{23}} C_7(m_W) + \frac{8}{3} \left( \eta^{\frac{14}{23}} - \eta^{\frac{16}{23}} \right) C_8(m_W) + \sum_{i=1}^{8} h_i \eta^{a_i}, \]

\[ C_8 = \eta^{\frac{14}{23}} C_8(m_W) + \sum_{i=1}^{8} \tilde{h}_i \eta^{\tilde{a}_i}, \]  

where $\eta = \alpha_s(m_W)/\alpha_s(m_b) \approx 0.56$, and $h_i, \tilde{h}_i$ and $a_i$ are known numerical coefficients. For the Wilson coefficients at the scale $m_W$, we write $C_{7,8}(m_W) = C_{7,8}(m_W) + C_{7,8}^{\text{new}}(m_W)$. The first term corresponds to the Standard Model contributions, which are functions of the mass ratio $x_t = (m_t/m_W)^2$. Numerically, one obtains

\[ C_7 \approx -0.31 + 0.67 C_7^{\text{new}}(m_W) + 0.09 C_8^{\text{new}}(m_W), \]

\[ C_8 \approx -0.15 + 0.70 C_8^{\text{new}}(m_W). \]  

(5)

We choose to parametrize our results in terms of the magnitude and phase of one of the New Physics contributions, $C_8^{\text{new}}(m_W) \equiv K_8 e^{i\gamma_8}$ or $C_7^{\text{new}}(m_W) \equiv -K_7 e^{i\gamma_7}$, as well as the ratio

\[ \xi = \frac{C_8^{\text{new}}(m_W)}{Q_d C_7^{\text{new}}(m_W)}, \]  

where $Q_d = -\frac{1}{3}$. A given New Physics scenario predicts these quantities at some large scale $M$. Using the renormalization group, it is then possible to evolve these predictions down to the scale $m_W$. Typically, $\xi \equiv \xi(M)$ tends to be smaller than $\xi(m_W)$ by an amount of order $-0.1$ to $-0.3$, depending on
Table 2: Ranges of $\xi(M)$ for various New Physics contributions to $C_7$ and $C_8$, characterized by the particles in penguin diagrams

| Class-1 models                                    | $\xi(M)$ |
|---------------------------------------------------|-----------|
| neutral scalar–vectorlike quark                   | 1         |
| gluino–squark ($m_{\tilde{g}} < 1.37m_{\tilde{q}}$) | $-0.13$–$-1$ |
| techniscalar                                      | $\approx -0.5$ |
| Class-2 models                                    | $\xi(M)$ |
| scalar diquark–top                                | 4.8–8.3   |
| gluino–squark ($m_{\tilde{g}} > 1.37m_{\tilde{q}}$) | $-1$–$-2.9$ |
| charged Higgs–top                                 | $-2.4$–$-3.8$ |
| left–right $W$–top                                | $\approx -6.7$ |
| Higgsino–stop                                     | $-2.6$–$-24$ |

how close the New Physics is to the electroweak scale. We restrict ourselves to cases where the parameter $\xi$ in (6) is real; otherwise there would be even more potential for CP violation. This happens if there is a single dominant New Physics contribution, such as the virtual exchange of a new heavy particle, contributing to both the magnetic and the chromo-magnetic dipole operators.

Ranges of $\xi(M)$ for several illustrative New Physics scenarios are collected in Table 2. For a detailed discussion of the model parameters which lead to the $\xi$ values quoted in the table the reader is referred to Ref. 6. Our aim is not to carry out a detailed study of each model, but to give an idea of the sizable variation that is possible in $\xi$. It is instructive to distinguish two classes of models: those with moderate (class-1) and those with large (class-2) values of $|\xi|$. It follows from (5) that for small positive values of $\xi$ it is possible to have large complex contributions to $C_8$ without affecting too much the magnitude and phase of $C_7$, since

$$
\frac{C_8}{C_7} \approx \frac{0.70K_8 e^{i\gamma_8} - 0.15}{(0.09 - 0.22\xi)K_8 e^{i\gamma_8} - 0.31}.
$$

This is also true for small negative values of $\xi$, albeit over a smaller region of parameter space. New Physics scenarios that have this property belong to class-1 and have been explored in Ref. 7. They allow for large CP asymmetries resulting from the $C_7$–$C_8$ interference term in (3). Figure 1 shows contour plots for the CP asymmetry in the $(K_8, \gamma_8)$ plane for six different choices of $\xi$ between $\frac{3}{2}$ and $-1$, assuming a cutoff $E_{\gamma} > 1.85\,\text{GeV}$ on the photon energy. We repeat that the results for the CP asymmetry depend very little on the
Figure 1: Contour plots for the CP asymmetry $A_{CP}^{b \rightarrow s\gamma}$ for various class-1 models. We show contours only until values $A_{CP} = 50\%$; for such large values, the theoretical expression for the CP asymmetry in (3) would have to be extended to higher orders to get a reliable result.

choice of the cutoff. For each value of $\xi$, the plots cover the region $0 \leq K_{8} \leq 2$ and $0 \leq \gamma_{8} \leq \pi$ (changing the sign of $\gamma_{8}$ would only change the sign of the CP asymmetry). The contour lines refer to values of the asymmetry of 1%, 5%, 10%, 15% etc. The thick dashed lines indicate contours where the $B \rightarrow X_{s}\gamma$ branching ratio takes values between $1 \times 10^{-4}$ and $4 \times 10^{-4}$, as indicated by the numbers inside the squares. The Standard Model prediction with this choice of the photon-energy cutoff is about $3 \times 10^{-4}$. The main conclusion to be drawn from Figure 1 is that in class-1 scenarios there exists great potential for sizable CP asymmetries in a large region of parameter space, whereas any point to the right of the 1% contour for $A_{CP}^{b \rightarrow s\gamma}$ cannot be accommodated by the Standard Model. Note that quite generally the regions of parameter space that yield large values for the CP asymmetries are not excluded by the experimental constraint on the CP-averaged branching ratio. To have large CP asymmetries the products $C_{i}C_{j}^{\ast}$ are required to have large imaginary parts [cf. (3)], whereas the total branching ratio is sensitive to the real parts of these quantities.
There are also scenarios in which the parameter $\xi$ takes on larger negative or positive values. In such cases, it is not possible to increase the magnitude of $C_8$ much over its Standard Model value, and the only way to get large CP asymmetries from the $C_7$–$C_8$ or $C_7$–$C_2$ interference terms in (3) is to have $C_7$ tuned to be very small; however, this possibility is constrained by the fact that the total $B \rightarrow X_s\gamma$ branching ratio must be of an acceptable magnitude. That this condition starts to become a limiting factor is already seen in the plots corresponding to $\xi = -\frac{1}{2}$ and $-1$ in Figure 1. For even larger values of $|\xi|$, the $C_7$–$C_8$ interference term becomes ineffective, because the weak phase tends to cancel in the ratio $C_8/C_7$. Then the $C_2$–$C_7$ interference term becomes the main source of CP violation; however, as discussed in Section 3, it cannot lead to asymmetries exceeding a level of about 5% without violating the constraint that the $B \rightarrow X_s\gamma$ branching ratio not be too small. Models of this type belong to the class-2 category. For a graphical analysis of class-2 models it is convenient to choose the magnitude and phase of the New Physics contribution $C_{new}^8(m_W) \equiv -K_7 e^{i\gamma_7}$ as parameters, rather than $K_8$ and $\gamma_8$. The reason is that for large $|\xi|$ it becomes increasingly unlikely that $C_{new}^8(m_W)$ will be large. The resulting plots are given in Figure 2. The branching-ratio constraint allows larger values of $C_8$ for positive $\xi$, which explains why larger asymmetries are attainable in this case. For example, for $\xi \approx 5$, which can be obtained from scalar diquark–top penguins, asymmetries of 5–20% are seen to be consistent with the $B \rightarrow X_s\gamma$ bound. On the other hand, for $\xi \approx -(2.5–5)$, which includes the multi-Higgs-doublet models, CP asymmetries of only a few percent are attainable, in agreement with the findings of previous authors. The same is true for the left–right symmetric $W$–top penguin, particularly if one takes into account that $K_7 \lesssim 0.2$ if $m_{W_R} > 1$ TeV.

The class-1 New Physics scenarios explored in Figure 1 have the attractive feature of a possible large enhancement of the magnitude of the Wilson coefficient $C_8$. This could have important implications for the phenomenology of the semileptonic branching ratio and charm yield in $B$ decays, through enhanced production of charmless hadronic final states induced by the $b \rightarrow sg$ transition. At $O(\alpha_s)$, the theoretical expression for the $B \rightarrow X_{sg}$ branching ratio is proportional to $|C_8|^2$. The left-hand plot in Figure 3 shows contours for this branching ratio in the $(K_8, \gamma_8)$ plane. In the Standard Model, $B(B \rightarrow X_{sg}) \approx 0.2\%$ is very small; however, in scenarios with $|C_8| = O(1)$ sizable values of order 10% for this branching ratio are possible, which simultaneously lowers the theoretical predictions for the semileptonic branching ratio and the charm production rate $n_c$ by a factor of $[1+B(B \rightarrow X_{sg})]^{-1}$. The most recent value of $n_c$ reported by the CLEO Collaboration is $1.12 \pm 0.05$. Although the systematic errors in this measurement are large, the result favours
Figure 2: Contour plots for the CP asymmetry $A_{CP}^{b\to s\gamma}$ for various class-2 models.

Figure 3: Contour plots for the $B \to X_{s\gamma}$ branching ratio (left) and for the charm yield $n_c$ in $B$ decays (right). There is an overall theoretical uncertainty of 6% on the values of $n_c$. 

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values of $B(B \to X_{s\gamma})$ of order 10%. This is apparent from the right-hand plot in Figure 3, which shows the central theoretical prediction for $n_c$ as a function of $K_8$ and $\gamma_8$. (There is an overall theoretical uncertainty in the value of $n_c$ of about 6% resulting from the dependence on quark masses and the renormalization scale.) The theoretical prediction for the semileptonic branching ratio would have the same dependence on $K_8$ and $\gamma_8$, with the normalization $B_{SL} = (12 \pm 1)%$ fixed at $K_8 = 0$. A large value of $B(B \to X_{s\gamma})$ could also help in understanding the $\eta'$ yields in charmless $B$ decays. For completeness, we note that the CLEO Collaboration has recently presented a preliminary upper limit on $B(B \to X_{s\gamma})$ of 6.8% (90% CL). It is therefore worth noting that large CP asymmetries of order 10–20% can also be attained at smaller $B \to X_{s\gamma}$ branching ratios of a few percent, which would nevertheless represent a marked departure from the Standard Model prediction.

4 Conclusions

I have reported on a study of direct CP violation in the inclusive, radiative decays $B \to X_{s\gamma}$. From a theoretical point of view, inclusive decay rates entail the advantage of being calculable in QCD, so that a reliable prediction for the CP asymmetry can be confronted with data. From a practical point of view, it is encouraging that $B \to X_{s\gamma}$ decays have already been observed experimentally, and high-statistics measurements will be possible in the near future. We find that in the Standard Model the CP asymmetry in $B \to X_{s\gamma}$ decays is strongly suppressed by three small parameters: $\alpha_s(m_b)$ arising from the necessity of having strong phases, $\sin^2\theta_C \approx 5\%$ reflecting a CKM suppression, and $(m_c/m_b)^2 \approx 8\%$ resulting from a GIM suppression. As a result, the asymmetry is only of order 1% in magnitude – a conclusion that cannot be significantly modified by long-distance contributions. We have argued that the latter two suppression factors are inoperative in extensions of the Standard Model for which the effective Wilson coefficients $C_7$ and $C_8$ receive additional contributions involving non-trivial weak phases. Much larger CP asymmetries are therefore possible in such cases.

A model-independent analysis of New Physics scenarios in terms of the magnitudes and phases of the Wilson coefficients $C_7$ and $C_8$ shows that, indeed, sizable CP asymmetries are predicted in large regions of parameter space. In particular, asymmetries of 10–50% are possible in models which allow for a strong enhancement of the coefficient of the chromo-magnetic dipole operator. They are, in fact, quite natural unless there is a symmetry that forbids new weak phases from entering the Wilson coefficients. Quite generally, having a large CP asymmetry is not in conflict with the observed value for the CP-
averaged \( B \to X_s\gamma \) branching ratio. On the contrary, it may even help to lower the theoretical prediction for this quantity, and likewise for the semileptonic branching ratio and charm multiplicity in \( B \) decays, thereby bringing these three observables closer to their experimental values.

The fact that a large inclusive CP asymmetry in \( B \to X_s\gamma \) decays is possible in many generic extensions of the Standard Model, and in a large region of parameter space, offers the possibility of looking for a signature of New Physics in these decays using data sets that will become available during the first period of operation of the \( B \) factories. A negative result of such a study would impose constraints on many New Physics scenarios. A positive signal, on the other hand, would provide interesting clues about the nature of physics beyond the Standard Model. In particular, a CP asymmetry exceeding the level of 10\% would be a strong hint towards enhanced chromo-magnetic dipole transitions caused by some new flavour physics.

We have restricted our analysis to the case of inclusive radiative decays since they entail the advantage of being very clean, in the sense that the strong-interaction phases relevant for direct CP violation can be reliably calculated. However, if there is New Physics that induces a large inclusive CP asymmetry in \( B \to X_s\gamma \) decays it will inevitably also lead to sizable asymmetries in some related processes. In particular, since we found that the inclusive CP asymmetry remains almost unaffected if a cut on the high-energy part of the photon energy spectrum is imposed, we expect that a large asymmetry will persist in the exclusive decay mode \( B \to K^*\gamma \), even though a reliable theoretical analysis would be much more difficult because of the necessity of calculating final-state rescattering phases.\footnote{Still, it would be worthwhile searching for a large CP asymmetry in this channel.}

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