SusyBSG: a fortran code for BR\([B \rightarrow X_s \gamma]\) in the MSSM with Minimal Flavor Violation

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Abstract

We present the fortran code SusyBSG version 1.1, which computes the branching ratio for the decay \(B \rightarrow X_s \gamma\) in the MSSM with Minimal Flavor Violation. The computation takes into account all the available NLO contributions, including the complete supersymmetric QCD corrections to the Wilson coefficients of the magnetic and chromomagnetic operators.

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1 Introduction

The radiative $B$ decays play a key role in the program of precision tests of the Standard Model (SM) and its extensions. The inclusive decay $B \to X_s \gamma$ is particularly well suited to this precision program, thanks to its low sensitivity to non-perturbative effects. The present experimental world average for the branching ratio of this decay, with a 1.6 GeV lower cut on the energy of the photon, is $\text{BR}[B \to X_s \gamma]_{\text{ex}} = (3.55 \pm 0.26) \times 10^{-4}$ [1]. The SM prediction for the branching ratio with the same cut on the photon energy is $\text{BR}[B \to X_s \gamma]_{\text{th}} = (3.15 \pm 0.23) \times 10^{-4}$ [2, 3] and includes most of the next-to-next-to-leading order (NNLO) perturbative QCD contributions as well as the leading non-perturbative and electroweak effects. Both experiment and SM prediction have an uncertainty of about 7%.

New Physics (NP) can in principle induce sizeable contributions to the decay $B \to X_s \gamma$, hence the good agreement between the SM prediction and the experimental result puts severe constraints on the flavor structure of NP models. However, the theoretical accuracy of the predictions for $\text{BR}[B \to X_s \gamma]$ in extensions of the SM is not at the same level as in the SM. A complete next-to-leading order (NLO) calculation is available only for the Two-Higgs-Doublet Model (THDM) [4, 5], the Left-Right symmetric model [6], and for the Minimal Supersymmetric Standard Model (MSSM) under the simplifying assumption of Minimal Flavor Violation (MFV), according to which the quark and squark mass matrices can be simultaneously diagonalized and the only source of flavor violation is the CKM matrix. In this scenario, the diagrams that include gluons and charginos were computed in refs. [7, 6], while the contributions involving gluinos were first considered in the heavy gluino limit in ref. [7], and in an effective Lagrangian approach in refs. [8, 9, 10]. After a partial two-loop calculation [11], the full computation of the two-loop gluino contributions to $B \to X_s \gamma$ in the MSSM with MFV was finally presented in ref. [12].

Several public computer codes that determine the MSSM mass spectrum and other SUSY observables (e.g. SuSpect [13], SPheno [14], micrOMEGAs [15], FeynHiggs [16], NMHDECAY [17], CPsuperH [18] and SuperIso [19]) contain calculations of $\text{BR}[B \to X_s \gamma]$ in various approximations. However, in the present versions of all these codes the two-loop gluino contributions to $B \to X_s \gamma$ are included, if at all, only in the effective Lagrangian approximation of refs. [8, 9, 10], which is valid in the limit of heavy superpartners and large $\tan \beta$.

In this paper we present a new fortran code, SuSyBSG, dedicated to the full NLO calculation of $\text{BR}[B \to X_s \gamma]$ in the MSSM with MFV. The code includes the full results of ref. [12] for the two-loop gluino contributions to the Wilson coefficients of the magnetic and chromomagnetic operators relevant to the $B \to X_s \gamma$ decay, and the results of refs. [7, 6] for the two-loop gluon contributions. It should be recalled that the weak interactions affect the squark and quark mass matrices in a different way, therefore their simultaneous diagonalization can be consistently imposed only at a scale $\mu_{\text{MFV}}$, which concurs to specify the MFV model. The renormalization group evolution of the MSSM parameters then leads to a disalignment between the squark and
quark mass matrices at scales different from $\mu_{\text{MFV}}$. Thus, optionally, the code allows for the inclusion of (small) additional contributions to the Wilson coefficients from one-loop diagrams with gluinos and down-type squarks (as well as charginos and up-type squarks) that occur when the MFV condition is imposed at a scale much higher than the weak scale. For the sake of comparison, SusyBSG can also provide evaluations of $\text{BR}[B \to X_s \gamma]$ in the SM and in the THDM, as well as in the MSSM with two-loop gluino contributions computed in the effective Lagrangian approximation.

In SusyBSG the relation between the Wilson coefficients and the $B \to X_s \gamma$ branching ratio is computed at NLO in perturbative QCD, along the lines of ref. [20], including also the dominant electroweak and non-perturbative corrections. For new physics that does not induce effective operators other than those already present in the SM, the NNLO anomalous dimensions of the effective operators and their matrix elements are the same as in the SM, suggesting the possibility of a partial NNLO implementation. A complete NNLO calculation would require also the NNLO contributions to the matching conditions, which have been computed only in the SM. However, if it can be argued that these contributions are as small as in the SM they can safely be neglected. This has been done for the case of the type II THDM in ref. [2], where the NP contributions to the matching conditions are computed at NLO and the anomalous dimensions and matrix elements are computed at NNLO. In the present version of SusyBSG we follow a similar but simpler route: we take into account the first NNLO estimate of ref. [2] by modifying the approach of ref. [20] in a way that approximately reproduces the results of a partial NNLO implementation. In fact, any attempt at a partial NNLO implementation has limitations in the MSSM, where higher-order QCD and electroweak contributions to the matching conditions can be sizeable, especially at large $\tan \beta$. The theoretical accuracy of our code therefore tends to be poorer in the MSSM than in the SM.

This manual is structured as follows: in section 2 we briefly summarize the two-loop results implemented in SusyBSG, focusing on the information necessary to a correct interpretation of the input parameters (we refer the readers to ref. [12] for the technical details on the calculation). In section 3 we describe the structure of SusyBSG, focusing in particular on the input and output parameters of the two main subroutines that make up the program. In section 4 we briefly detail our default choices for the input parameters and discuss the theoretical uncertainty. In the appendices we provide additional details on the various corrections implemented in SusyBSG. In the appendix A we present the one-loop gluino contributions to the matching conditions for the Wilson coefficients $C_1^{(1)}$ and $C_2^{(1)}$. In the appendix B we provide formulae for the one-loop gluino and chargino contributions to the Wilson coefficients in the presence of flavor mixing in the squark sector. In the appendix C we summarize our treatment of the $\tan \beta$-enhanced contributions to the Wilson coefficients. Finally, in the appendix D we summarize the NLO computation of the relation between Wilson coefficients and branching ratio for $B \to X_s \gamma$.

The latest version of SusyBSG can be downloaded from the Web page

http://cern.ch/slavich/susybsg/home.html

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2 Radiative $B$ decays in the MSSM with MFV

In this section we summarize the calculation of the weak-scale matching conditions for the $\Delta B = 1$ effective Hamiltonian in the MSSM with Minimal Flavor Violation, as implemented in SusyBSG. The NLO relation between the Wilson coefficients at the weak scale and the branching ratio is summarized for completeness in the appendix D.

The $\Delta B = 1$ effective Hamiltonian at the matching scale $\mu_0$ (of the order of the weak scale) is given by

$$H = -\frac{4G_F}{\sqrt{2}} V_{ts}^{\text{CKM}} V_{tb}^{\text{CKM}} \sum_i C_i(\mu_0) Q_i(\mu_0),$$

where $G_F$ is the Fermi constant and $V_{ts}^{\text{CKM}}$, $V_{tb}^{\text{CKM}}$ are elements of the CKM matrix. The operators relevant to our calculation are

$$\begin{align*}
Q_1 &= (\bar{s}_L \gamma_\mu T^a c_L)(\bar{c}_L \gamma^\mu T^a b_L), \\
Q_2 &= (\bar{s}_L \gamma_\mu c_L)(\bar{c}_L \gamma^\mu b_L), \\
Q_3 &= (\bar{s}_L \gamma_\mu b_L) \sum_q (\bar{q} \gamma_\mu q), \\
Q_4 &= (\bar{s}_L \gamma_\mu T^a b_L) \sum_q (\bar{q} \gamma_\mu T^a q), \\
Q_5 &= (\bar{s}_L \gamma_\mu \gamma_\nu \gamma_\rho b_L) \sum_q (\bar{q} \gamma_\mu \gamma_\nu \gamma_\rho q), \\
Q_6 &= (\bar{s}_L \gamma_\mu \gamma_\nu \gamma_\rho T^a b_L) \sum_q (\bar{q} \gamma_\mu \gamma_\nu \gamma_\rho T^a q), \\
Q_7 &= \frac{e}{16\pi^2} m_b \bar{s}_L \sigma^{\mu \nu} b_R F_{\mu \nu}, \\
Q_8 &= \frac{g_s}{16\pi^2} m_b \bar{s}_L \sigma^{\mu \nu} T^a b_R G_{\mu \nu}^a.
\end{align*}$$

When the QCD corrections are considered, the Wilson coefficients of the operators $Q_i$ can be organized in the following way

$$C_i(\mu_0) = C_i^{(0)\text{SM}}(\mu_0) + C_i^{(0)H^\pm}(\mu_0) + C_i^{(0)\text{SUSY}}(\mu_0) + \frac{\alpha_s(\mu_0)}{4\pi} \left[ C_i^{(1)\text{SM}}(\mu_0) + C_i^{(1)H^\pm}(\mu_0) + C_i^{(1)\text{SUSY}}(\mu_0) \right],$$

where the various leading order (LO) contributions are classified according to whether the corresponding diagrams contain only SM fields, a physical charged Higgs boson and an up-type quark, or a chargino and an up-type squark. The expressions for $C_i^{(0)\text{SM}}$ and $C_i^{(0)H^\pm}$ can be found, e.g., in ref. [4], while those for $C_i^{(0)\text{SUSY}}$ can be found, e.g., in eq. (4) of ref. [7]. Note that at LO only the coefficients of the magnetic and chromo-magnetic operators $Q_7$ and $Q_8$ receive contributions from non-SM fields. One-loop neutralino- and gluino-exchange diagrams should be neglected under the MFV assumption.
The NLO coefficients $C^{(1)}_{\text{SM}}$ and $C^{(1)}_{i}H^{\pm}$ contain the gluonic corrections to the SM and charged Higgs contributions, respectively, and can be found for instance in ref. [4]. Concerning the NLO supersymmetric contributions $C^{(1)}_{i}^{\text{SUSY}}$, the chargino-gluon contributions can be found in refs. [7, 6], and a full two-loop computation of the gluino contributions to $C^{(1)}_{7,8}^{\text{SUSY}}$ in the MSSM with minimal flavor violation was more recently presented in ref. [12]. Together with the one-loop gluino contributions to $C^{(1)}_{1,2}^{\text{SUSY}}$ that we present in the appendix A, the results of ref. [12] provide us with a complete NLO computation of the supersymmetric QCD contributions to $\text{BR}[B \to X_s \gamma]$ in the MFV scenario.

The study of NLO contributions in the MFV scenario is complicated by the fact that the simultaneous diagonalization of squark and quark mass matrices is not preserved by radiative corrections. As a result, even when the renormalized mixing matrices for quarks and squarks are assumed to be flavor-diagonal, flavor-changing counterterms of electroweak origin have to be taken into account in the two-loop calculation. The technical issues related to the renormalization of flavor mixing are discussed in ref. [12], to which we refer the interested reader. In the following we summarize the information necessary to a correct interpretation of the input parameters required by SusyBSG.

In the MFV framework the mass matrices for both the up-type and down-type quarks and squarks are assumed to be simultaneously flavor-diagonal at some renormalization scale $\mu_{\text{MFV}}$. When computing the matching conditions for the Wilson coefficients we neglect the masses of the first- and second-generation quarks, as well as the left-right mixing in the mass matrices of the corresponding squarks. Thus, the soft SUSY-breaking terms that enter the squark mass matrices and are relevant to our calculation are: the masses for the SU(2) squark doublets, $m_{Q_i}$, where $i$ is a generation index; the masses for the third-generation singlets, $m_T$ and $m_B$; the trilinear interaction terms for the third-generation squarks, $A_t$ and $A_b$. We recall that in the so-called super-CKM basis, where the matrices of Yukawa couplings are diagonal and the squarks are rotated parallel to the quarks, the $3 \times 3$ mass matrices for the up-type and down-type left squarks are related by $(M_U^0)_{LL} = V^{\text{CKM}} (M_D^0)_{LL} V^{\text{CKM}}$. Therefore, the two mass matrices can be both flavor-diagonal only if they are flavor-degenerate. This means that the MFV scenario can be consistently implemented only if we choose a common mass parameter for the three generations of SU(2) squark doublets, i.e. $m_{Q_i} \equiv m_Q$.

Since we are focusing on the QCD corrections at two loops, it is necessary to specify a renormalization scheme for the input parameters $m_Q$, $m_T$ and $A_t$, which determine the up-type squark masses and mixing entering $C^{(0)}_{7,8}^{\text{SUSY}}$ and are subject to $O(\alpha_s)$ radiative corrections. We consider two options: the first is to assume that they are Lagrangian parameters expressed in a minimal subtraction scheme such as $\overline{\text{DR}}$, at some renormalization scale $\mu_{\text{SUSY}}$ of the order of the superparticle masses. The second is to adopt an on-shell (OS) definition: $m_Q$, $m_T$ and $A_t$ can be interpreted as the unphysical parameters that enter the tree-level stop mass matrix obtained by rotating the diagonal matrix of the physical stop masses by a suitably defined physical mixing angle $\theta_t$. In this case we adopt an OS definition (i.e. we use the physical mass)
also for the top quark mass that enters the tree-level stop mass matrix.

The other MSSM parameters relevant to the calculation, for which we need not specify a renormalization prescription, are: the ratio of Higgs vacuum expectation values \( \tan \beta \equiv v_2/v_1 \); the charged Higgs boson mass \( m_{H^\pm} \); the gluino mass \( m_{\tilde{g}} \); the SU(2) gaugino mass parameter \( M_2 \); the higgsino mass parameter \( \mu \). Our conventions for the signs of the various Lagrangian parameters are as specified by the SUSY Les Houches Accord (SLHA) [21, 22], i.e.: the top and bottom quark masses are
\[
m_t = h_t v_2 / \sqrt{2} \quad \text{and} \quad m_b = h_b v_1 / \sqrt{2},
\]
respectively; the left-right mixing terms in the stop and sbottom mass matrices are
\[
m_t(A_t - \mu \cot \beta) \quad \text{and} \quad m_b(A_b - \mu \tan \beta),
\]
respectively; the (2,2) entry of the chargino mass matrix is \( \mu \).

SusyBSG has an option to read the parameters of the MSSM Lagrangian from a SLHA spectrum file. In that case, the parameters are meant to be expressed in the \( \overline{\text{DR}} \) renormalization scheme. When a SLHA spectrum file is available the program reads from it also the \( \overline{\text{DR}} \)-renormalized top quark mass computed at the scale \( \mu_{\text{SUSY}} \), which is necessary for the computation of the running stop masses. In the absence of a SLHA spectrum file the running top mass is computed internally by the program, taking into account only the corrections controlled by the strong gauge coupling.

In a two-loop calculation the interpretation of the MFV requirement itself depends on the way we renormalize the flavor mixing, i.e. the way we fix the counterterms that cancel the divergences of the antihermitian parts of the quark and squark wave-function-renormalization (WFR) matrices. In particular, if we perform a minimal subtraction we are imposing the MFV condition on the \( \overline{\text{DR}} \)-renormalized parameters of the Lagrangian evaluated at the scale \( \mu_{\text{MFV}} \). In this scheme \( C^{(1)_{\text{SUSY}}} \) contains logarithms of the ratio \( M_S/\mu_{\text{MFV}} \), where \( M_S \) represents the mass of the superparticles entering the loops. An alternative option consists in subtracting also the finite part of the antihermitian WFR: this results in a conventional (and gauge-dependent) on-shell renormalization scheme [23], in which \( C^{(1)_{\text{SUSY}}} \) is independent of \( \mu_{\text{MFV}} \).

From the discussion above it is clear that, in models where the MFV condition is imposed at a scale much larger than the superparticle masses (such as, e.g., supergravity models where one identifies \( \mu_{\text{MFV}} \) with the GUT scale), the Wilson coefficients computed in the minimal subtraction scheme contain very large logarithms of \( M_S/\mu_{\text{MFV}} \). In this case, no matter what renormalization scheme is chosen, the fixed-order calculation does not provide a good approximation to the correct result. Indeed, in such models the soft SUSY-breaking mass parameters – which are flavor-diagonal at the scale \( \mu_{\text{MFV}} \) – must be evolved down to \( M_S \) with the appropriate RGE, thus generating some flavor violation in the squark mass matrices. When the squark mass matrices are diagonalized, the resummed logarithms of the ratio \( M_S/\mu_{\text{MFV}} \) are absorbed in the couplings of the resulting squark mass eigenstates with the gluinos (and the charginos and neutralinos). Typically, the effects of the RGE-induced flavor mixing are relatively small, and we include them only at leading order\(^1\) in the one-loop diagrams with gluinos and down-type quarks.

\(^1\)We follow the approach of ref. [24]. For a more complete treatment of the gluino contributions in scenarios with generic sources of flavor violation see ref. [25].
squarks (which would vanish if the MFV condition was valid at the low scale $M_S$) and in the one-loop diagrams with charginos and up-type squarks. The corresponding contributions to the Wilson coefficients are given explicitly in the appendix B. Once the logarithmic effects have been taken into account in this way, the genuine two-loop MFV contributions in $C^{(1)\text{SUSY}}_{7,8}$ can be computed by setting artificially $\mu_{\text{MFV}} \sim M_S$. In this case, using either the on-shell scheme or the minimal subtraction scheme to renormalize the flavor mixing will give basically the same result.

We remark here that the rather complicated task of solving the system of RGE equations for the soft SUSY-breaking parameters, taking into account the full flavor structure of the MSSM, is not performed by SusyBSG. In order to obtain the squark masses and mixing matrices that enter the expressions in appendix B, starting from flavor-universal boundary conditions at the scale $\mu_{\text{MFV}}$, the users can run one of the public spectrum calculators that include a full treatment the flavor structure, such as the latest (still unpublished) versions of SPheno \cite{hahn2014} or SoftSusy \cite{heinemeyer2013}, and pass the results to SusyBSG by means of a SLHA2 \cite{heinemeyer2013} input/output file.

Finally, it is well known that in the MSSM the relation between the bottom quark mass $m_b$ and the bottom Yukawa coupling $h_b$ is subject to tan $\beta$-enhanced threshold corrections \cite{hurth2014}. If the bottom Yukawa coupling entering the one-loop part of the Wilson coefficients is expressed in terms of the SM value of the running bottom mass, the SUSY threshold corrections induce counterterm contributions that, although being formally of higher order in the loop expansion, are enhanced by powers of tan $\beta$ and may therefore be sizeable when tan $\beta$ is large. As discussed e.g. in ref. \cite{hurth2014}, such potentially large corrections can be absorbed in the one-loop results by a suitable redefinition of the bottom Yukawa coupling. Other tan $\beta$-enhanced contributions to the Wilson coefficients appear in the form of corrections to the Higgs-quark-quark vertices, and have been computed in refs. \cite{heinemeyer2013} in an effective Lagrangian approach where the heavy SUSY particles are integrated out of the theory. Among the tan $\beta$-enhanced two-loop contributions, those controlled by the strong gauge coupling are fully accounted for by the calculation of ref. \cite{heinemeyer2013}. In addition, we include in SusyBSG the tan $\beta$-enhanced contributions controlled by the top and bottom Yukawa couplings, following the approach of ref. \cite{heinemeyer2013}. More details on the treatment of the tan $\beta$-enhanced contributions are given in the appendix C.
3 Structure of the program SusyBSG

SusyBSG 1.1 is structured in two modules, which in principle could be used independently as stand-alone programs. The first module is the subroutine `WilsonCoeff`, which computes the weak-scale matching conditions for the Wilson coefficients of the $\Delta B = 1$ effective Hamiltonian in a physics model to be chosen among the SM, the Two Higgs Doublet Model and the MSSM. The second module is the subroutine `getBR`, which computes $\text{BR}[B \to X_s \gamma]$ at NLO, taking as input the SM and new-physics contributions to the Wilson coefficients evaluated by `WilsonCoeff`. The main program `BSGAMMA`, which can be freely modified according to the users’ needs, sets all the relevant input parameters, calls the two subroutines and prints out the output. Within `BSGAMMA` the users are allowed to read the input parameters required by `WilsonCoeff` from a spectrum file written in the SLHA format. More specifically, the subroutine `readSUSY_SLHA1` reads from a SLHA1 [21] file only the flavor-conserving parameters required for the calculation of the one- and two-loop contributions appearing in eq. (10). The subroutine `readSUSY_SLHA2`, instead, reads from a SLHA2 [22] file all the parameters of the MSSM Lagrangian, including the flavor-violating parameters necessary to compute the additional one-loop contributions discussed in the appendix B. In either case the spectrum file must be named `SLHA.in`, and it must be located in the directory where the program is run.

For certain choices of the input parameters it may happen that the masses of two particles are accidentally very similar to each other, or that the sum of two masses is very close to a third mass. In some pathologic cases this leads to numerical instabilities in the output of SusyBSG, even if all the formulae for the two-loop gluino contributions to $B \to X_s \gamma$ are well behaved when the corresponding limits are taken analytically. Since those formulae are very long and depend in a complicated way on the values of twelve different particle masses, providing analytical results to cover all the problematic limits would be highly inefficient in terms of size and speed of the code. Therefore, we limit ourselves to looking for the occurrence of accidentally similar masses, and issuing a warning if necessary. We leave it to the users to check that, in those cases, the result for $\text{BR}[B \to X_s \gamma]$ is not unreasonably sensitive to small variations of the input parameters.\footnote{Note that the numerical instabilities can be greatly reduced by compiling the program in quadruple precision. Details are provided in the SusyBSG Web page.}

The only exception is the case in which the mass of the first- or second-generation up-type squark $m_{\tilde{u}_L}$ is very similar to the mass of the super-strange $m_{\tilde{s}_L}$. This happens inevitably if the soft SUSY-breaking parameter $m_Q$, common to both masses, is much larger than $m_Z$. To avoid numerical instabilities we switch to the analytical results valid in the limit $m_{\tilde{u}_L} = m_{\tilde{s}_L}$ if the relative difference between the two masses is less than 1%.

In the following we describe in detail the input and output parameters of the two main subroutines that make up SusyBSG.
3.1 The subroutine WilsonCoeff

The call to the subroutine for the matching conditions to the Wilson coefficients reads

```fortran
    call WilsonCoeff(imod,scheme,mu0,mususy,mumfv,
        $  msq3,mstr,msbr,msql,At,Ab,mHp,mg,M2,mu,tanb,
        $  CISM,C7SM,C8SM,CINP,C7NP,C8NP,prob,eqmass)
```

The variables in the first two lines of the `call` command are inputs, and are defined as:

- **integer imod**: allows the users to choose the particle content of the theory and the approximation used in the computation of the two-loop contribution to the Wilson coefficients. The values 0–4 correspond to:
  - 0: Standard Model fields only. The new-physics contributions are set to zero;
  - 1: Two Higgs Doublet Model. The only additional contributions come from diagrams with a charged Higgs $H^\pm$;
  - 2: MSSM with the effective Lagrangian approach. At NLO, only the $\tan \beta$-enhanced SUSY contributions are included in the Wilson coefficients as in ref. [10];
  - 3: MSSM with full NLO QCD contributions. Includes the two-loop calculation of the gluino contributions presented in ref. [12], and the results of ref. [10] for the remaining (i.e., non-QCD) $\tan \beta$-enhanced contributions.
  - 4: MSSM with full NLO QCD contributions (as for imod = 3), with in addition the contributions of the one-loop diagrams with gluinos and down-type squarks and chargino and up-type squarks given in the appendix B.

- **logical scheme(2)**: contains two logical switches (relevant only for imod = 2,3,4) that allow the users to specify the renormalization conditions in the squark sector. In particular:
  - 1: choice of renormalization scheme for the input squark mass parameters
    
    ```fortran
    (.true. = OS, .false. = DR);
    ```
  - 2: choice of renormalization scheme for the MFV condition
    
    ```fortran
    (.true. = OS, .false. = DR).
    ```

  If the MSSM Lagrangian parameters are read in from a SLHA file both entries of `scheme` should be set to `.false.`.

- **real*8 mu0**: renormalization scale $\mu_0$ (of the order of the weak scale) at which the matching of the Wilson coefficients is performed.
• real*8 mususy: renormalization scale $\mu_{\text{SUSY}}$ at which the input squark mass parameters are given (relevant only for scheme(1) = .false.).

• real*8 mumfv: renormalization scale $\mu_{\text{MFV}}$ at which the MFV condition is imposed (relevant only for scheme(2) = .false.). As explained in section 2, the results of the program are not reliable if $\mu_{\text{MFV}}$ is set to be much larger than $\mu_{\text{SUSY}}$.

• real*8 msq3,mstr,msbr: third-generation squark mass parameters $m_{Q_3}$, $m_T$ and $m_B$.

• real*8 msql: mass parameter $m_Q$ for the first- and second-generation squark doublets. Note that the MFV condition is only consistent with $m_Q = m_{Q_3}$.

• real*8 At,Ab: third-generation Higgs-squark-squark interaction terms $A_t$ and $A_b$.

• real*8 mHp: mass of the charged Higgs boson $m_{H^\pm}$.

• real*8 mg: gluino mass $m_{\tilde{g}}$.

• real*8 M2: SU(2) gaugino mass parameter (enters the chargino masses).

• real*8 mu: Higgs-mixing superpotential parameter $\mu$.

• real*8 tanb: ratio of Higgs vacuum expectation values $\tan \beta$.

of course, all of the SUSY input parameters are relevant only for imod = 2,3,4, while mHp and tanb are relevant for imod = 1,2,3,4. The option imod = 4 requires the presence of a SLHA2 spectrum file, from which the subroutine readSUSY_SLHA2, to be called before WilsonCoeff, reads the SUSY input parameters. For consistency, when imod = 4 the entries of scheme should be both .false., and $\mu_{\text{MFV}}$ should be set close to $\mu_{\text{SUSY}}$.

The variables in the third line of the call command are outputs, and are defined as:

• real*8 CISM(6): vector containing the Standard Model contributions to $C_i^{(1)}(\mu_0)$ for $1 \leq i \leq 6$ (they are actually different from zero only for $i = 1, 4$).

• real*8 C7SM(2),C8SM(2): vectors whose two elements are the Standard Model contributions to $C_{7,8}^{(0)}(\mu_0)$ and $C_{7,8}^{(1)}(\mu_0)$.

• real*8 CINP(6): vector containing the new-physics contributions to $C_i^{(1)}(\mu_0)$ for $1 \leq i \leq 6$ (in the MSSM they are different from zero only for $i = 1, 2, 4$).

• real*8 C7NP(2),C8NP(2): vectors whose two elements are the new-physics contributions to $C_{7,8}^{(0)}(\mu_0)$ and $C_{7,8}^{(1)}(\mu_0)$. 


• **logical prob**: problem flag, set to `.true.` if the relative difference between two (or more) particle masses is less than 1%. The case $m_{\tilde{u}_L} \approx m_{\tilde{s}_L}$ – which is taken care of internally – is not flagged, nor are some other cases that never lead to instabilities.

• **logical eqmass(12)**: vector that identifies the masses that are too close to each other. The ordering is $(m_t, m_W, m_\tilde{q}, m_{H^\pm}, m_{\tilde{b}_1}, m_{\tilde{b}_2}, m_{\tilde{t}_1}, m_{\tilde{t}_2}, m_{\tilde{u}_L}, m_{\tilde{s}_L}, m_{\chi^+_1}, m_{\chi^+_2})$, and the entries of `eqmass` corresponding to the masses that are too close are set to `.true.`

In addition to the input variables passed in the call to the subroutine, the users must set (or read from a SLHA spectrum file) a number of Standard Model parameters in the common block `SMINPUTS`:

```plaintext```
common/SMINPUTS/mz,mw,mtpole,mbmb,hsm,asmz,azinv
```

The variables in `SMINPUTS` are defined as:

- **real*8 mz,mw**: physical masses $m_Z$ and $m_W$ for the SM gauge bosons.
- **real*8 mtpole**: physical top-quark mass $m_t^{\text{pole}}$.
- **real*8 mbmb**: running bottom mass $m_b^{\overline{\text{MS}}}(m_b)$ in the $\overline{\text{MS}}$ scheme at the scale $m_b$.
- **real*8 hsm**: mass $m_h$ of the SM-like Higgs boson (used only for the small electroweak corrections of ref. [29], see the appendix D).
- **real*8 asmz,azinv**: strong and inverse electromagnetic couplings $\alpha_s(m_Z)$ and $1/\alpha(m_Z)$ in the $\overline{\text{MS}}$ scheme at the scale $m_Z$.

### 3.2 The subroutine `getBR`

This subroutine computes $\text{BR}[B \to X_s \gamma]$ at NLO using the results of refs. [20, 29]. For the users’ convenience the computation of the branching ratio will be summarized in the appendix D. The call to the subroutine reads:

```plaintext```
call getBR(muw,mut,mub,muc,E0,CINP,C7NP,C8NP,BR)
```

The variable `BR` is the output value for $\text{BR}[B \to X_s \gamma]$. All the other variables in the call are inputs, and are defined as:

- **real*8 muw**: weak scale $\mu_w \sim m_W$ at which the light quark contributions to the Wilson coefficients are computed.
- **real*8 mut**: weak scale $\mu_t \sim m_t$ at which the top quark contributions to the Wilson coefficients are computed. The new-physics contributions (see below) should also be computed by the subroutine `WilsonCoeff` at a scale $\mu_0 = \mu_t$. 

• **real*8 mub**: low scale $\mu_b \sim m_b$ at which the branching ratio is computed.

• **real*8 muc**: low scale $\mu_c$ at which the charm quark mass $m_c^{\overline{\text{MS}}}(\mu_c)$ entering the SM contributions to the Wilson coefficients is computed.

• **real*8 E0**: minimum photon energy $E_0$.

• **real*8 CINP(6)**: vector containing the new-physics contributions to $C_i^{(1)}(\mu_t)$ for $1 \leq i \leq 6$ (they are produced in output by the subroutine WilsonCoeff).

• **real*8 C7NP(2),C8NP(2)**: vectors whose two elements are the new-physics contributions to $C_7^{(0)}(\mu_t)$ and $C_8^{(1)}(\mu_t)$ (they are produced in output by the subroutine WilsonCoeff).

In addition to the input variables passed in the call, the subroutine getBR requires the Standard Model input parameters contained in the common block SMINPUTS (see section 3.1), and a set of parameters contained in the common block BRINPUTS

```
common/BRINPUTS/a0inv,mcmc,rbs,hlam,ccsl,bsl,lambda,A,rhobar,etabar
```

The variables in BRINPUTS are defined as

• **real*8 a0inv**: inverse of the fine-structure constant, $1/\alpha_{\text{em}}$.  

• **real*8 mcmc**: running charm mass $m_c^{\overline{\text{MS}}}(m_c)$ the $\overline{\text{MS}}$ scheme at the scale $m_c$.  

• **real*8 rbs**: ratio $m_b/m_s$ entering the gluon bremsstrahlung contribution.  

• **real*8 hlam**: HQET parameter $\lambda_2$ entering the non-perturbative contribution.  

• **real*8 ccsl**: non-perturbative semileptonic phase-space factor $C$.  

• **real*8 bsl**: semileptonic branching ratio $\text{BR}[B \rightarrow X_c e \bar{\nu}]$.  

• **real*8 lambda,A,rhobar,etabar**: Wolfenstein parameters $\lambda$, $A$, $\bar{\rho}$ and $\bar{\eta}$ for the CKM matrix.

More information on these quantities can be found in the appendix D. Note that the parameters entering BRINPUTS are *not* read in from a SLHA spectrum file (even when one is present) and must be explicitly set by the users before calling getBR.
3.3 File structure of SusyBSG

The latest version of SusyBSG can be downloaded from the program’s Web page in the form of a file named SusyBSG_vv.tar.gz, where vv stands for the version number. When uncompressed and unpacked, the program consists of the following files:

- **BSGAMMA.f**: contains the main program BSGAMMA, to be modified (or replaced) by the users. It just sets the relevant inputs, calls the two subroutines WilsonCoeff and getBR and prints out the output.

- **DGStwoloop.f**: contains the subroutine DGSbsgamma, which computes the two-loop gluino contributions to the Wilson coefficients, based on the results of ref. [12]. Upon compilation it includes all the files located in the directory source.

- **functions.f**: contains the definitions of various functions entering the one-loop and two-loop contributions to the Wilson coefficients.

- **getBR.f**: contains the subroutine getBR, based mostly on the results of ref. [20], which computes the branching ratio for the process $B \rightarrow X_s \gamma$ taking as input the new-physics contributions to the Wilson coefficients.

- **Makefile**: by default it uses the compiler g77 to produce an executable named SusyBSG. Further instructions for the compilation can be found in the program’s Web page.

- **slha2io.f**: contains two routines that read the input parameters (with the exception of those in the common block BRINPUTS) from a spectrum file in the SLHA format [21, 22]. readSUSY_SLHA1 reads only the parameters necessary to the calculation of the MFV contributions, while readSUSY_SLHA2 reads also the parameters necessary to the calculation of the one-loop diagrams that involve flavor-violating gluino-quark-squark or chargino-quark-squark vertices.

- **SLHA.in**: an example of SLHA1 spectrum file that can be used as input by SusyBSG. The file was produced with SoftSusy, with input parameters corresponding to the so-called SPS1a' [30] point.

- **source**: the files in this directory contain the explicit expressions of the two-loop gluino contributions to the Wilson coefficients, to be included in DGStwoloop.f upon compilation. The files were automatically converted to fortran from Mathematica format using FormCalc [31].

- **WilsonCoeff.f**: contains the subroutine WilsonCoeff, which computes the one-loop and two-loop new-physics contributions to the Wilson coefficients (calling DGSbsgamma for the two-loop gluino contributions).
4 Default values of the SM input parameters

In a phenomenological study of $BR[B \to X_s \gamma]$ in the MSSM with MFV, the SM input parameters entering the common blocks $\text{SMINPUTS}$ and $\text{BRINPUTS}$ will presumably be fixed once and for all at the beginning of the calculation. For the users’ convenience we list below the default values of these parameters in the version 1.1 of SusyBSG. In most cases we use the same values as in ref. [3], where the users can look for the corresponding references.

The default values of the SM parameters entering the common block $\text{SMINPUTS}$ are:

\begin{align*}
\text{mz} &= m_Z = 91.1876 \text{ GeV} \\
\text{mw} &= m_W = 80.403 \text{ GeV} \\
\text{mtpole} &= m_t^{\text{pole}} = 170.9 \text{ GeV} \\
\text{mbls} &= m_b^{\text{MS}}(m_b) = 4.2 \text{ GeV} \\
\text{hsm} &= m_h = 115 \text{ GeV} \\
\text{asmz} &= \alpha_s(m_Z) = 0.1189 \\
\text{azinv} &= 1/\alpha(m_Z) = 127.918
\end{align*}

The default values of the B-physics parameters entering the common block $\text{BRINPUTS}$ are:

\begin{align*}
\text{a0inv} &= 1/\alpha_{\text{em}} = 137.036 \\
\text{mcmc} &= m_c^{\text{MS}}(m_c) = 1.224 \text{ GeV} \\
\text{rbs} &= m_b/m_s = 50 \\
\text{hlam} &= \lambda_2 = 0.12 \text{ GeV}^2 \\
\text{ccsl} &= C = 0.580 \\
\text{bsl} &= \text{BR}[B \to X_c e \bar{\nu}]_{\text{exp}} = 0.1061 \\
\text{lambda} &= \lambda = 0.2272 \\
\text{A} &= A = 0.818 \\
\text{rhobar} &= \bar{\rho} = 0.221 \\
\text{etabar} &= \bar{\eta} = 0.340
\end{align*}

The choice of the four independent renormalization scales that appear in the calculation of the branching ratio and are required as input by the subroutine $\text{getBR}$ deserves a separate discussion. We recall that these scales are: the weak scales $\mu_t$ and $\mu_W$, at which we compute the contributions to the matching conditions for the Wilson coefficients coming from top and charm quarks, respectively; the low scale $\mu_b$ at which we compute the matrix elements for the $b \to s\gamma$ decay; the scale $\mu_c$ at which we express the charm mass entering the matrix elements. We adopt the default values

$$\mu_t = \mu_W = 2 m_W, \quad \mu_b = 2.5 \text{ GeV}, \quad \mu_c = 1 \text{ GeV}. \quad (11)$$

Using these values for the scales, and setting $m_t^{\text{pole}} = 171.4 \text{ GeV}, \text{mb}_s^{\text{MS}}(m_b) = 4.15 \text{ GeV}$ and $1/\alpha(m_Z) = 128.940$ in order to reproduce the inputs of ref. [2], we obtain a SM prediction
for \( \text{BR}[B \to X_s \gamma] \) of 3.15 \( \times 10^{-4} \), in agreement with the NNLO result of ref. [2]. Very good agreement, within up to 2\%, is also found at various values of \( \tan \beta \) and \( m_{H^\pm} \) with the partial NNLO implementation of the THDM from refs. [2,32]. While we plan to include the known NNLO corrections in a future version of SusyBSG, the present implementation provides an excellent starting point.

It should be noted, however, that \( \mu_c \) in eq. (11) is adjusted to a very low value in order to mimic the NNLO contributions that are not present in our calculation. Therefore, in this case, the variation of the renormalization scales should not be used to estimate the intrinsic uncertainty of our calculation. Indeed, the result of the NLO calculation depends quite sharply on \( \mu_c \) around the value that reproduces the NNLO result. For example, using \( \mu_c = m_{c}^{\text{MS}}(m_c) = 1.224 \text{ GeV} \), with the other scales fixed as in eq. (11), results in \( \text{BR}[B \to X_s \gamma] = 3.29 \times 10^{-4} \).

Concerning the theoretical uncertainty of our prediction, we recall that in the SM analysis of refs. [2,3] the error is dominated by a 5\% uncertainty due to unknown \( O(\alpha_s \Lambda_{\text{QCD}}/m_b) \) non-perturbative contribution to the matrix elements. Additional \( \sim 4\% \) intrinsic uncertainty stems from the perturbative part of the calculation and from the estimate of missing NNLO contributions. The parametric error in the SM is only about 3\%. All these errors are present in the MSSM calculation as well. Therefore, we recommend using at least a 7\% intrinsic uncertainty, throughout the parameter space, to be added in quadrature with the parametric uncertainty. Barring the case of subtle cancellations between various supersymmetric contributions, this appears to be a realistic guesstimate of our theory uncertainty.

To conclude, we quote the results of SusyBSG 1.1 in a characteristic point of the MSSM parameter space, the so called SPS1a′ point [30]. We used SoftSusy to generate a SLHA1 spectrum file that provides the subroutine WilsonCoeff with the DR-renormalized SUSY parameters evaluated at a scale \( \mu_{\text{SUSY}} \) equal to the geometric average of the stop masses. We set \( \mu_{\text{MFV}} \) to be equal to \( \mu_{\text{SUSY}} \), thus enforcing the MFV condition at the weak scale. We also set the SM input parameters to the default values listed above (with the exception of \( m_W \) and \( m_h \), which are taken from the output of SoftSusy), and the renormalization scales appearing in the calculation of the branching ratio to the values given in eq. (11). For the SM prediction (obtained with \( \text{imod} = 0 \)) we find that \( \text{BR}[B \to X_s \gamma] = 3.16 \times 10^{-4} \). For the THDM prediction (\( \text{imod} = 1 \)) we find 3.91 \( \times 10^{-4} \). For the MSSM prediction in the effective Lagrangian approximation (\( \text{imod} = 2 \)) we find 2.42 \( \times 10^{-4} \). Finally, for the MSSM prediction based on the full NLO QCD calculation (\( \text{imod} = 3 \)) we find 2.60 \( \times 10^{-4} \).

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Appendix A: gluino contributions to $C^{(1)}_1$ and $C^{(1)}_2$

A complete NLO computation of the supersymmetric QCD contributions to $\text{BR}[B \to X_s \gamma]$ requires the one-loop contributions of $\mathcal{O}(\alpha_s)$ to the Wilson coefficients of the four-fermion operators $Q_1-6$. All the contributions from Standard Model particles, as well as those arising from loops involving a charged Higgs boson and a quark or a chargino and a squark, can be found in the literature \[4, 7, 6\]. There are however additional contributions to the Wilson coefficients $C^{(1)}_1$ and $C^{(1)}_2$ that originate from loops involving gluinos, see figs. \[1\] and \[2\]. These contributions were neglected in earlier computations performed under the assumption of heavy gluinos, but they must be taken into account in a general computation. Assuming MFV, and neglecting $m_c$ and $m_s$ as well as the left-right mixing in the charm- and strange-squark sectors, the SUSY contributions to $C^{(1)}_1$ and $C^{(1)}_2$ at the matching scale $\mu_0$ read

\[
C^{(1)\text{SUSY}}_1(\mu_0) = \frac{m_{\chi_1} m_{\tilde{g}_2}}{m_g^2} \left[ X_{ij}^b L_{ij}^b X_{i} L_1^s (x_{\chi_i}, x_{\tilde{b}_j}, x_{\tilde{c}_L}) + L_i^c L_i^s (x_{\chi_i}, x_{\tilde{s}_L}, x_{\tilde{c}_L}) \right] + \frac{m_{\tilde{g}_2}^2}{2 m_g^2} \left[ |L_i^s|^2 I_2 (x_{\chi_i}, x_{\tilde{c}_L}) + X_{ij}^b L_{ij}^b I_2 (x_{\chi_i}, x_{\tilde{b}_j}, x_{\tilde{c}_L}) \right],
\]

\[
C^{(1)\text{SUSY}}_2(\mu_0) = -\frac{C_F}{2} \left[ |X_{i1}^b|^2 I_3 (x_{\tilde{b}_i}, x_{\tilde{c}_L}) + I_3 (x_{\tilde{s}_L}, x_{\tilde{c}_L}) \right],
\]

where: summation over the repeated indices is understood; $X^b$ is the mixing matrix for the sbottoms, defined as $(\tilde{b}_1, \tilde{b}_2)^T = X^b (\tilde{b}_L, \tilde{b}_R)^T$; $C_F = 4/3$ is a color factor; the quark-squark-chargino couplings are defined as

\[
L_{ij}^b = \frac{m_b}{m_{\tilde{g}} \cos \beta} U_{i2} X_{j2}^b - \sqrt{2} U_{i1} X_{j1}^b, \quad L_i^s = -\sqrt{2} U_{i1}, \quad L_i^c = -\sqrt{2} V_{i1}^*,
\]

where $U$ and $V$ are the unitary matrices that diagonalize the chargino mass matrix according to $U \mathcal{M}_\chi V^T = \text{diag}(m_{\chi_1}, m_{\chi_2})$. Finally, $x_P \equiv m_P^2/m_g^2$ for a generic particle $P$, and the loop integrals $I_i$ are defined as

\[
I_1(x_1, x_2, x_3) = \frac{x_1 \ln x_1}{(1-x_1)(x_1-x_2)(x_2-x_3)} + (x_1 \leftrightarrow x_2) + (x_1 \leftrightarrow x_3),
\]

\[
I_2(x_1, x_2, x_3) = \frac{-x_1^2 \ln x_1}{(1-x_1)(x_1-x_2)(x_1-x_3)} + (x_1 \leftrightarrow x_2) + (x_1 \leftrightarrow x_3),
\]

\[
I_3(x_1, x_2) = \frac{x_1 + x_2 - 2 x_1 x_2}{2 (x_1 - 1)(x_2 - 1)} + \frac{x_1 (x_1^2 - 2 x_2 x_1 + x_1 x_2) \ln x_1}{(x_1 - 1)^2 (x_1 - x_2)} + (x_1 \leftrightarrow x_2).
\]
Figure 1: Box diagrams contributing to the Wilson coefficient $C^{(1)_{\text{SUSY}}}_{1}$.

Figure 2: Vertex diagrams contributing to the Wilson coefficient $C^{(1)_{\text{SUSY}}}_{2}$. Self-energy insertions with a squark-gluino loop on the external quark legs are also taken into account.
Appendix B: contributions from squark flavor mixing

In this appendix we summarize the well-known results for the gluino–squark and chargino–squark contributions to the Wilson coefficients in presence of flavor mixing in the squark sector. In the super-CKM basis, where the squark fields are rotated parallel to their fermionic partners, the mass eigenstates are linear combinations of flavor eigenstates, according to

\[
\begin{pmatrix}
\tilde{u}_1 \\
\tilde{u}_2 \\
\tilde{u}_3 \\
\tilde{u}_4 \\
\tilde{u}_5 \\
\tilde{u}_6
\end{pmatrix} = \begin{pmatrix}
\Gamma_{UL} & \Gamma_{UR}
\end{pmatrix}
\begin{pmatrix}
\tilde{u}_L \\
\tilde{c}_L \\
\tilde{t}_L \\
\tilde{u}_R \\
\tilde{c}_R \\
\tilde{t}_R
\end{pmatrix},
\begin{pmatrix}
\tilde{d}_1 \\
\tilde{d}_2 \\
\tilde{d}_3 \\
\tilde{d}_4 \\
\tilde{d}_5 \\
\tilde{d}_6
\end{pmatrix} = \begin{pmatrix}
\Gamma_{DL} & \Gamma_{DR}
\end{pmatrix}
\begin{pmatrix}
\tilde{d}_L \\
\tilde{s}_L \\
\tilde{t}_L \\
\tilde{b}_L \\
\tilde{s}_R \\
\tilde{b}_R
\end{pmatrix},
\tag{B1}
\end{equation}

where \( \Gamma_{QL} \) and \( \Gamma_{QR} \) (with \( Q=U,D \)) are 6 \times 3 mixing matrices. The contributions to the Wilson coefficients from one-loop diagrams with gluinos and down-type squarks, \( C^{(0)}_{7,8} \), have to be added to the first line in eq. (10). They read:

\[
C^{(0)}_{7,8} = \frac{g_s^2 m_w^2}{g^2 m_{\tilde{g}}^2} V_{ts}^{CKM} V_{tb}^{CKM} \sum_i \left[ (\Gamma^*_{DL})_{i2} (\Gamma_{DL})_{i3} f_{7,8} \left( \frac{m_{\tilde{d}_i}^2}{m_{\tilde{g}}^2} \right) - \frac{m_{\tilde{g}}}{m_b} (\Gamma^*_{DL})_{i2} (\Gamma_{DR})_{i3} g_{7,8} \left( \frac{m_{\tilde{d}_i}^2}{m_{\tilde{g}}^2} \right) \right],
\tag{B2}
\]

where \( g_s \) and \( g \) are the strong and SU(2) gauge couplings, respectively, and the functions \( f_{7,8} \) and \( g_{7,8} \) are defined as:

\[
f_7(x) = \frac{2(2 + 5x - x^2)}{27(x-1)^3} - \frac{4x}{9(x-1)^2} \ln x, \quad g_7(x) = -\frac{4(1 + x)}{9(x-1)^2} + \frac{8x}{9(x-1)^3} \ln x, \tag{B3}
\]

\[
f_8(x) = \frac{(11 - 40x + 19x^2)}{36(x-1)^3} - \frac{x(1 - 9x)}{6(x-1)^4} \ln x, \quad g_8(x) = -\frac{5 - 13x}{3(x-1)^2} + \frac{x(1 - 9x)}{3(x-1)^3} \ln x. \tag{B4}
\]

When the information about the mixing in the squark sector is available we replace the chargino-squark contributions to \( C^{(0)\text{SUSY}}_{7,8} \) appearing in eq. (10) with

\[
C^{(0)\text{SUSY}}_{7,8} = \sum_{i,j,l,j} \frac{V_{is}^{CKM} V_{js}^{CKM} V_{is}^{CKM}}{V_{ts}^{CKM} V_{tb}^{CKM}} \left[-L^{I*}_{ij} L^{I}_{ij} \frac{m_w^2}{3 m_{\tilde{u}_j}^2} F^3_{7,8} \left( \frac{m_{\tilde{u}_j}^2}{m_{\tilde{g}}^2} \right) + L^{I*}_{ij} R^{I}_{ij} \frac{m_w^2}{2 m_{\tilde{g}}^2} F^3_{7,8} \left( \frac{m_{\tilde{u}_j}^2}{m_{\tilde{g}}^2} \right) \right],
\tag{B5}
\]

where: the indices \( i,j \) run over the up-type squark flavors; the functions \( F^3_{7,8} \) can be read, e.g., in eqs. (2.4) and (3.2)–(3.3) of ref. [8]; the chargino-squark-squark couplings \( L^{I}_{ij} \) and \( R^{I}_{ij} \) read

\[
L^{I}_{ij} = \frac{m_{\tilde{u}}^I m_u^I}{m_w \sin \beta} V_{i2}^* (\Gamma_{UR})_{j1} - \sqrt{2} V_{i1}^* (\Gamma_{UL})_{j1}, \quad R^{I}_{ij} = \frac{1}{\cos \beta} U_{i2}^* (\Gamma_{UL})_{j1}, \tag{B6}
\]

where the matrices \( U \) and \( V \) are defined after eq. (A3). We checked that eqs. (B2) and (B5) agree with the corresponding results in ref. [24].
Appendix C: \( \tan \beta \)-enhanced contributions

In this appendix we describe in detail the treatment of the \( \tan \beta \)-enhanced contributions to the Wilson coefficients relevant to \( b \to s\gamma \). The contributions arising from the threshold corrections to the relation between the bottom mass and the bottom Yukawa coupling \([27]\) can be absorbed in the one-loop part of the Wilson coefficients by a suitable redefinition of the Yukawa coupling \( h_b \). Indicating by \(-\epsilon_b\) the \( \tan \beta \)-enhanced part of the SUSY contribution to the bottom quark self-energy and by \(-\delta_b\) the rest of the SUSY contribution, so that the relation between the running bottom masses in the SM and in the MSSM is

\[
m_{b}^{\text{SM}} = m_{b}^{\text{MSSM}} (1 + \epsilon_b \tan \beta + \delta_b),
\]

we multiply the contributions of the one-loop diagrams that involve a bottom Yukawa coupling by a factor \( \kappa \), defined as

\[
\kappa = \frac{1 - \delta_b}{1 + \epsilon_b \tan \beta}.
\]

This amounts to defining \( h_b \) in terms of the MSSM running bottom mass in the one-loop contributions to the Wilson coefficients, with the result that the counterterm contributions at higher orders do not contain terms enhanced by powers of \( \tan \beta \). For consistency we multiply by \( \kappa \) also the contributions of the two-loop diagrams that involve a bottom Yukawa coupling, but we neglect there the small effect of \( \delta_b \).

The \( \mathcal{O}(\alpha_s) \) contributions to \( \epsilon_b \) and \( \delta_b \) read

\[
\epsilon_b^{(s)} = -\frac{\alpha_s(\mu_0)}{3\pi} \frac{2\mu}{m_{\tilde{g}}} H_2 \left( \frac{m_{\tilde{b}_1}^2}{m_{\tilde{g}}^2}, \frac{m_{\tilde{b}_2}^2}{m_{\tilde{g}}^2} \right),
\]

\[
\delta_b^{(s)} = \frac{\alpha_s(\mu_0)}{3\pi} \left\{ \frac{2 A_b}{m_{\tilde{g}}} H_2 \left( \frac{m_{\tilde{b}_1}^2}{m_{\tilde{g}}^2}, \frac{m_{\tilde{b}_2}^2}{m_{\tilde{g}}^2} \right) - \left[ B_1(0, m_{\tilde{g}}^2, m_{\tilde{b}_1}^2, \mu_0^2) + B_1(0, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2, \mu_0^2) \right] \right\},
\]

where \( H_2(x, y) \) is defined in eq. (2.8) of ref. [8], and \( B_1(p^2, m_1^2, m_2^2, \mu^2) \) is a Passarino-Veltman function defined as in ref. [34].

In addition to the \( \tan \beta \)-enhanced contributions that arise from the threshold corrections to the bottom mass, there are \( \tan \beta \)-enhanced contributions arising from corrections to the Higgs-quark-quark vertices \([8, 9, 10]\). We stress that all the contributions of this kind controlled by the strong gauge coupling are fully accounted for by the explicit two-loop calculation of ref. [12].

The \( \tan \beta \)-enhanced contributions to the Wilson coefficients controlled by the top and bottom Yukawa couplings are implemented in SusyBSG using the effective Lagrangian approach of refs. [8, 9, 10]. They include a contribution of \( \mathcal{O}(\alpha_t) \) to the bottom quark self-energy

\[
\epsilon_b^{(t)} = -\frac{\alpha_t(\mu_0)}{4\pi} \frac{A_t}{\mu} H_2 \left( \frac{m_Q^2}{\mu^2}, \frac{m_T^2}{\mu^2} \right),
\]
such that we must use $\epsilon_b = \epsilon_b^{(s)} + \epsilon_b^{(t)}$ when computing the factor $\kappa$ in eq. (C2). In addition, there is a contribution to the Wilson coefficients arising from a correction of $\mathcal{O}(\alpha_t)$ to the effective $G^+-t-b$ vertex:

$$\delta C_{7,8} = -\frac{\epsilon_b^{(t)} \tan \beta}{1 + \epsilon_b \tan \beta} F_{7,8}^{(2)} \left( \frac{m_t^2}{m_H^2} \right),$$

where the loop functions $F_{7,8}^{(2)}(x)$ are defined e.g. in eq. (2.4) of ref. [8]. In the limit of heavy superparticles the contribution in eq. (C6) cancels an analogous term originating from the redeffinition of the bottom Yukawa coupling, thus ensuring the decoupling of new-physics effects from the SM contribution.

There are also contributions to the Wilson coefficients involving the bottom Yukawa coupling and higher powers of $\tan \beta$:

$$\delta C_{7,8} = -\frac{\epsilon_b^{(t)} \epsilon_t^{(b)} \tan^2 \beta}{(1 + \epsilon_b \tan \beta)(1 + \epsilon_b^{(s)} \tan \beta)} F_{7,8}^{(2)} \left( \frac{m_t^2}{m_H^2} \right)$$

$$-\frac{\epsilon_b^{(t)} \tan \beta}{(1 + \epsilon_b \tan \beta)^2(1 + \epsilon_b^{(s)} \tan \beta)} \frac{m_b^2}{36 m_H^2},$$

where the $\mathcal{O}(\alpha_b)$ vertex correction $\epsilon_t^{(b)}$ is defined as

$$\epsilon_t^{(b)} = \frac{\alpha_b(\mu_0)}{4\pi} \frac{A_b}{\mu} H_2 \left( \frac{m_Q^2}{\mu^2}, \frac{m_B^2}{\mu^2} \right).$$

Being suppressed by an additional loop factor and by $m_b^2/m_H^2$, respectively, the two terms in eq. (C7) tend to be numerically small, but, as stressed in ref. [10], they can become relevant in special cases where the leading $\mathcal{O}(\alpha_s)$ and $\mathcal{O}(\alpha_t)$ effects cancel each other.

When implementing in SusyBSG the tan $\beta$-enhanced contributions controlled by the top and bottom Yukawa coupling, eqs. (C5)–(C8), we have neglected the mixing between superpartners, approximating the squark masses with the corresponding soft SUSY-breaking terms and the masses of the higgsino components of charginos and neutralinos with the superpotential parameter $\mu$. Indeed, the effective Lagrangian approach used to derive these results relies on the assumption that the superpartners are much heavier than the weak scale and can be integrated out of the theory, leaving behind only non-decoupling corrections to the Higgs-quark-quark vertices. In this case the effect of the mixing, which is due to the breaking of electroweak symmetry, can reasonably be expected to be small, and anyway is of the same order of magnitude as other effects that are neglected under the effective Lagrangian approximation$^3$.

$^3$Formulæ for the tan $\beta$-enhanced contributions to the Wilson coefficients that take into account the effect of superpartner mixing have been presented in refs. [8] [15] [35] [36].
Appendix D: NLO determination of $\text{BR}[B \to X_s \gamma]$

In this appendix we summarize the NLO computation of the relation between the branching ratio for $B \to X_s \gamma$ and the Wilson coefficients computed at the electroweak scale, as implemented in SusyBSG. We rely mainly on the results of ref. [20], but we also take into account the matrix elements for the four-fermion operators computed in ref. [37] (resulting in an updated table of “magic numbers”) and the electroweak matching conditions computed in ref. [29].

The $B \to X_s \gamma$ branching ratio with a photon energy cutoff $E_0$ in the $B$-meson rest frame can be related to the experimentally measured rate $\text{BR}[B \to X_c e \bar{\nu}]_{\text{exp}}$ as

$$\frac{\text{BR}[B \to X_s \gamma]_{E_0 > E_0}}{\text{BR}[B \to X_c e \bar{\nu}]_{\text{exp}}} = \left| \frac{V_{ts}^{\text{CKM}} V_{tb}^{\text{CKM}}}{V_{cb}^{\text{CKM}}} \right|^2 \frac{6 \alpha_{\text{em}}}{\pi} C \left( |K|^2 + B(E_0) + N(E_0) \right),$$

(D1)

where $C$ is the “non-perturbative semileptonic phase-space factor” defined in the appendix C of ref. [20]. The quantities $B(E_0)$ and $N(E_0)$ represent the gluon-bremsstrahlung and non-perturbative contributions, respectively, and will be discussed below. The factor $K$ contains the contributions to the $b \to s \gamma$ amplitude and is dominated by the effective Wilson coefficient for the magnetic operator at the low scale. It can be written as

$$K = K_c + r(\mu_t) \left[ K_t + K_{\text{NP}} \right] + \varepsilon_{\text{ew}}.$$

(D2)

Following ref. [20] we separate $K$ into the light quark contribution $K_c$, the top-quark contribution $K_t$, the new-physics contribution $K_{\text{NP}}$ and the electroweak correction $\varepsilon_{\text{ew}}$. In the top and NP contributions we keep the bottom Yukawa coupling renormalized at the matching scale $\mu_t$ by introducing the quantity $r(\mu_t)$, defined as

$$r(\mu_t) \equiv \frac{m_b^{\text{MS}}(\mu_t)}{m_b^{\text{MS}}(\mu_b)} = \frac{m_b^{\text{MS}}(\mu_t)}{m_b^{\text{MS}}(\mu_b)} \times \left[ 1 - \frac{\alpha_s(m_b^{1S})}{4\pi} \left( \frac{16}{3} + 8 \ln \frac{m_b^{1S}}{\mu_b} \right) + \frac{2}{9} \alpha_s(m_b^{1S})^2 \right],$$

(D3)

where the “1S mass” $m_b^{1S}$ is defined as half of the perturbative contribution to the $\Upsilon$-mass. We extract the value of $m_b^{1S}$ from the input value of $m_b^{\text{MS}}(m_b)$ using eq. (168) of ref. [38]. Unlike ref. [20], we evaluate $r(\mu_t)$ strictly at NLO. Indeed, in eq. (D3) the ratio $r(\mu_t)$ is expressed as the product of a term that accounts for the evolution of the running bottom mass from $\mu_t$ to $\mu_b$ and a term (in the square brackets) that accounts for the shift between $m_b^{\text{MS}}(\mu_b)$ and $m_b^{1S}$. The NLO relation between the values of a running quark mass evaluated at two different renormalization scales $\mu$ and $\mu_0$ reads:

$$m_q^{\text{MS}}(\mu) = m_q^{\text{MS}}(\mu_0) \left[ \frac{\alpha_s(\mu)}{\alpha_s(\mu_0)} \right]^{\beta_0/\beta_1} \left[ 1 + \frac{\alpha_s(\mu_0)}{4\pi} \frac{\gamma_0}{\beta_0} \left( \frac{\gamma_1 - \beta_1}{\gamma_0 - \beta_0} \right) \left( \frac{\alpha_s(\mu)}{\alpha_s(\mu_0)} - 1 \right) \right],$$

(D4)

where $\beta_0 = 11 - \frac{2}{3} n_f$, $\beta_1 = 102 - \frac{38}{3} n_f$, $\gamma_0 = 8$ and $\gamma_1 = \frac{404}{3} - \frac{40}{3} n_f$ ($n_f$ being the number of active quark flavors). The strong gauge coupling $\alpha_s(\mu)$ is extracted at NLO from $\alpha_s(m_Z)$
according to
\[
\alpha_s(\mu) = \frac{\alpha_s(m_Z)}{v} \left( 1 - \beta_1 \frac{\alpha_s(m_Z) \ln v}{\beta_0 \frac{4\pi}{v}} \right), \quad v = 1 + \frac{\beta_0 \alpha_s(m_Z)}{2\pi} \ln \frac{\mu}{m_Z}.
\]  \quad (D5)

The light-quark contribution to $K$ reads
\[
K_c = -\frac{23}{36} \frac{\eta_\tau^3}{\eta_\tau^2} - \frac{8}{9} \left( \frac{14}{\eta_\tau^2} - \frac{16}{\eta_\tau^2} \right) + \sum_{k=1}^{8} h_k \eta_{W_k}^{a_k} + \frac{\alpha_s(\mu_b)}{4\pi} \sum_{k=1}^{8} \eta_{W_k}^{a_k} \left[ \frac{46}{3} a_k d_k \left( \ln \frac{m_1^{1S}}{\mu_b} + \eta_w \ln \frac{\mu_w}{m_w} \right) + \tilde{d}_k + \tilde{d}_k^0 \eta_w + \tilde{d}_k^0 a(z) + \tilde{d}_k^0 b(z) \right]
\]
\[
+ \frac{\sqrt{V_{ts}^{\text{CKM}} V_{tb}^{\text{CKM}}}}{V_{ts}^{\text{CKM}}} \frac{\alpha_s(\mu_b)}{4\pi} \left( \eta_{W_3}^{a_3} + \eta_{W_4}^{a_4} \right) \left[ a(z) + b(z) \right], \quad (D6)
\]
where: $\mu_w$ is the matching scale (not necessarily equal to $\mu_t$); $\eta_w = \alpha_s(\mu_w)/\alpha_s(\mu_b)$; the “magic numbers” $h_k$ are given in eq. (2.3) of ref. [20]; the “magic numbers” $a_k, d_k, \tilde{d}_k^0$ and $\tilde{d}_k^0$ are given in table 2 of ref. [37]; the functions $a(z)$ and $b(z)$ are given in eqs. (D.1) and (D.2) of ref. [20]. The variable $z$ is defined as $(m_c^{\text{MS}}/m_b^{1S})^2$, where $m_c$ is an arbitrary renormalization scale.

After factoring out $r(\mu_t)$, the top-quark contribution to $K$ reads
\[
K_t = \left[ 1 - \frac{2}{9} \alpha_s(m_b^{1S})^2 \right] \left[ \bar{\eta}_t^2 C_7^{(0)\text{SM}}(\mu_t) + \frac{8}{3} \left( \eta_t^2 - \frac{8}{9} \right) C_8^{(0)\text{SM}}(\mu_t) \right]
\]
\[
+ \frac{\alpha_s(\mu_b)}{4\pi} \left\{ \sum_{k=1}^{8} e_k \eta_t^{(a_k+\frac{1}{3})} C_4^{(1)\text{SM}}(\mu_t) \right\}
\]
\[
+ \eta_t^2 \left[ \eta_t C_7^{(1)\text{SM}}(\mu_t) - 2 \left( \frac{12523}{3174} - \frac{7411}{4761} \eta_t - \frac{2}{9} \eta_t^2 - \frac{4}{3} \ln \frac{m_b^{1S}}{\mu_b} \right) C_7^{(0)\text{SM}}(\mu_t) \right]
\]
\[
- \frac{8}{3} \eta_t C_8^{(1)\text{SM}}(\mu_t) - 2 \left( \frac{50092}{357075} + \frac{1110842}{357075} \eta_t + \frac{16}{27} \eta_t^2 + \frac{32}{9} \ln \frac{m_b^{1S}}{\mu_b} \right) C_8^{(0)\text{SM}}(\mu_t) \right]\]
\[
+ \eta_t^2 \left[ \frac{8}{3} \eta_t C_8^{(1)\text{SM}}(\mu_t) - 2 \left( \frac{2745458}{357075} - \frac{38890}{14283} \eta_t - \frac{4}{9} \eta_t^2 - \frac{16}{9} \ln \frac{m_b^{1S}}{\mu_b} \right) C_8^{(0)\text{SM}}(\mu_t) \right],
\]  \quad (D7)

where: $\eta_t = \alpha_s(\mu_t)/\alpha_s(\mu_b)$; the “magic numbers” $e_k$ are given in table 2 of ref. [37]; the top quark contributions to the matching conditions for the Wilson coefficients at the scale $\mu_t$ can be expressed in terms of the functions $E_0^t, A_{0,1}^t$ and $F_{0,1}^t$ appearing in eqs. (6), (10)–(12) and (17) of ref. [39] as
\[
C_4^{(1)\text{SM}}(\mu_t) = E_0^t(x), \quad C_7^{(0,1)\text{SM}}(\mu_t) = -\frac{1}{2} A_{0,1}^t(x), \quad C_8^{(0,1)\text{SM}}(\mu_t) = -\frac{1}{2} F_{0,1}^t(x),
\]  \quad (D8)
where \( x = (m_t^{\overline{MS}}(\mu_t)/m_W)^2 \). To compute \( x \), the \( \overline{\text{MS}} \)-renormalized running top mass is extracted from the pole mass \( m_t^{\text{pole}} \) according to

\[
m_t^{\overline{\text{MS}}}(m_t^{\text{pole}}) = m_t^{\text{pole}} \left[ 1 + \frac{4}{3} \frac{\alpha_s(m_t^{\text{pole}})}{\pi} + 10.9 \frac{\alpha_s^2(m_t^{\text{pole}})}{\pi^2} \right]^{-1},
\]

and is then evolved to the renormalization scale \( \mu_t \) according to eq. (D4).

The new-physics contribution to \( K \) can be obtained by simply replacing the top contributions in eq. (D7) with the corresponding NP contributions. In the MSSM one has to add also the squark-gluino contributions to the Wilson coefficients \( C_1 \) and \( C_2 \), presented in the appendix A. In summary,

\[
K_{NP} = K_t \left( C_n^{(i)\text{SM}} \to C_n^{(i)\text{NP}} \right) + \frac{\alpha_s(\mu_b)}{4\pi} \sum_{k=1}^{8} \left[ f_k C_1^{(1)\text{NP}}(\mu_t) + h_k C_2^{(1)\text{NP}}(\mu_t) \right] \eta_t^{(a_k+\frac{11}{2})},
\]

where the “magic numbers” \( h_k \) are the same as in \( K_c \), and the \( f_k \) are:

\[
f_k = (0.5784, -0.3921, -0.1429, 0.0476, -0.1275, 0.0317, 0.0078, -0.0031). \tag{D11}
\]

In the determination of the electroweak contribution \( \varepsilon_{ew} \) that appears in eq. (D2) we adopt the results of ref. [29] [see in particular eq. (3.10) there], including the available NP contributions where appropriate.

The gluon-bremsstrahlung contribution to the branching ratio is

\[
B(E_0) = \frac{\alpha_s(\mu_b)}{\pi} \sum_{i,j=1,\ldots,8}^{i,j \leq 8} C_i^{(0)\text{eff}}(\mu_b) C_j^{(0)\text{eff}}(\mu_b) \phi_{ij}(\delta), \tag{D12}
\]

where: \( \delta \equiv 1 - 2E_0/m_b^{1S} \); the expression for the functions \( \phi_{ij}(\delta) \) are given in eqs. (E.2)–(E.8) of ref. [20]; the ratio \( m_b/m_s \) appearing in the expression of \( \phi_{88}(\delta) \) is taken as input. The contributions corresponding to \( 3 \leq i \leq 6 \) or \( 3 \leq j \leq 6 \) are negligible and we set them to zero. We take the SM contributions to the LO effective Wilson coefficients \( C_i^{(0)\text{eff}}(\mu_b) \) from eq. (E.9) of ref. [20], and include the NP contributions where appropriate.

Finally, the non-perturbative contribution to the branching ratio can be approximated as (see [20] for the complete list of references)

\[
N(E_0) \simeq - \frac{1}{18} C_7^{(0)\text{eff}}(\mu_b) \left[ \frac{\eta_{w}}{\eta_{w} + \eta_{\nu}} \right] \frac{\lambda_2}{m_c^2}, \tag{D13}
\]

Where the parameter \( \lambda_2 \) is taken as input, and for the charm quark mass we use \( m_{c}^{\overline{\text{MS}}}(m_c) \).
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PROGRAM SUMMARY

Manuscript Title: SusyBSG: a fortran code for BR[B → Xsγ] in the MSSM with Minimal Flavor Violation

Authors: G. Degrassi, P. Gambino, P. Slavich

Program Title: SusyBSG

Journal Reference:  

Catalogue identifier: None

Licensing provisions: None

Programming language: Fortran

Operating system: Linux

Keywords: Supersymmetry, B physics, rare decays

PACS: 12.60.Jv, 13.20.He

Classification: 11.6 Phenomenological and Empirical Models and Theories

Nature of problem:
Predicting the branching ratio for the decay B → Xs γ in the MSSM with Minimal Flavor Violation.

Solution method:
We take into account all the available NLO contributions, including the two-loop gluino contributions to the Wilson coefficients computed in ref. [1]. The relation between the Wilson coefficients and the branching ratio for B → Xs γ is computed along the lines of ref. [2]. The SUSY input parameters can be read from a spectrum file in the SLHA format.

Restrictions:
The results apply only to the case of Minimal Flavor Violation.

Unusual features:
Numerical instabilities may arise for special combinations of the input parameters. They can usually be avoided by compiling the code in quadruple precision.

Running time:
Approximately 10 ms in double precision, 0.5 s in quadruple precision (on a Pentium D 3.40-GHz).

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[1] G. Degrassi, P. Gambino and P. Slavich, Phys. Lett. B 635 (2006) 335 [arXiv:hep-ph/0601135].
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