The rare $\bar{B}_d^0 \rightarrow \phi \gamma$ decays in Standard Model and as a probe of R-parity violation

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Abstract

We present the first study of the rare annihilation decay $\bar{B}_d^0 \rightarrow \phi \gamma$ in the Standard Model. Using QCD factorization formalism, we find $B(\bar{B}_d^0 \rightarrow \phi \gamma) = 3.6 \times 10^{-12}$. The smallness of the decay rate in the Standard Model make the decay a sensitive probe of new physics contributions. As an example, we calculate the effects of R-parity violating couplings. Within the available upper bounds for $|\lambda_{i23}^{''}\lambda_{i12}''|$ and $|\lambda_{i32}\lambda_{i12}^{'*}|$, $B(\bar{B}_d^0 \rightarrow \phi \gamma)$ could be enhanced to order of $10^{-9} \sim 10^{-8}$, which might be accessible at LHCb, B-TeV and the planning super high luminosity B factories at KEK and SLAC.

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1 Introduction

It is known that flavor changing neutral currents (FCNC) induced rare B decays are very sensitive probes of new physics. The GIM suppression of FCNC amplitude is absent in many new physics scenarios beyond the Standard Model (SM), which could give large enhancement of FCNC processes over the SM predictions. To search such kind of signals is one of the most important goals of B projects BaBar, Belle, BTeV, and LHC-B. On the other hand, B rare decays also serve as laboratory for hadronic dynamics. Due to our poor knowledge for non-perturbative QCD, predictions for many interesting decays are always polluted by uncomfortable large uncertainties, which have hindered us very much in extracting weak interaction information precisely from the available measurements. The well known example is two-body charmless B decays. It would be of great interesting to explore B rare decays which are induced by FCNC currents as well as involve few hadronic parameters.

To the end, we will study the pure penguin annihilation decay $\bar{B}_d^0 \to \phi \gamma$. Experimentally, $\bar{B}_d^0 \to \phi \gamma$ is very easy to be identified through the decay chain $\bar{B}_d^0 \to \phi \gamma \to (K^-K^+)\gamma$, i.e, two charge tracks and one energetic photon, and the detecting efficiency should be high. However, if the decay rate is too rare, it would be very difficult to pick out the signals of the decay from its background, which comes from continuum $(e^+e^- \to q\bar{q}$ with $q=u,d,s$) events with high energy photons originating from initial state radiation or $e^+e^- \to (\pi^0\eta)X$ with $\pi^0\eta \to \gamma\gamma$[1]. CLEO Collaboration had performed a search for the decay, but found no evidence and put the upper limit $\mathcal{B}(\bar{B}_d^0 \to \phi \gamma) < 0.33 \times 10^{-5}$ at 90% CL. To our best knowledge, there is no realistic theoretical study of $\bar{B}_d^0 \to \phi \gamma$ in the literature. In this paper, we will investigate this decay. In naive factorization approach, this decay involve the simple matrix $\langle \phi | \bar{s} \gamma_\mu | 0 \rangle$ and the same hadronic matrix element $\langle \gamma | \bar{q} \gamma_\mu (1 - \gamma_5) b | \bar{B}_q \rangle$ as the radiative leptonic decay which has been studied in Ref.[2, 3, 4, 5] with different framework. Beyond naive factorization, non-factorizable contribution should be included. The QCD factorization framework[6], which has been developed recently by Beneke et al., is employed to calculate $\mathcal{O}(\alpha_s)$ nonfactorization contributions arising from exchanging hard gluon between the two color octet currents $\bar{s}_\alpha \gamma_\mu (1 \pm \gamma_5) s_\beta$ and $\bar{d}_\beta \gamma_\mu (1 - \gamma_5) b_\alpha$. We find $\mathcal{B}(\bar{B}_d^0 \to \phi \gamma) = 3.6 \times 10^{-12}$ in the SM. The decay rate is too small to be observed at the running B factories BaBar and BELLE. Any measurement of the decay at BaBar and BELLE would be the evidence of activity of New Physics.
example, we treat it as a probe of R parity violating couplings (RPV). Within the available lowest upper bound of RPV couplings, it is found that the branching ratio of the decay could be enhanced to $10^{-9} \sim 10^{-8}$, which might be measured at LHC-B, BTeV and the planning super high luminosity B factories at KEK and SLAC.

2 $\bar{B}_d^0 \to \phi \gamma$ in the Standard Model

The $\bar{B}_d^0 \to \phi \gamma$ decay is illustrated schematically in fig.1. which is dominated by the photon radiating from the light quark in the B meson. Obviously, the amplitude of fig.1 is suppressed by power of $\Lambda_{QCD}/m_b$ because $\phi$ meson must be transversely polarized and B meson is heavy. The diagrams with the photon radiating from the heavy b quark and the energetic strange quarks of the $\phi$ meson are suppressed by power of $\Lambda_{QCD}^2/m_b^2$, which will be neglected in this paper. The situation is similar to that of the annihilation contributions in $B \to K^*\gamma$ decays\(^7\).

![Figure 1](image)

Figure 1: The leading diagram for $\bar{B}_d^0 \to \phi \gamma$. Radiation of the photon from the remaining three quark lines are further suppressed by power of $\Lambda_{QCD}/m_b$. The symbol $\otimes$ denotes the insertion of penguin operators $O_{3-10}$.

We begin our study with the effective Hamiltonian in the SM relevant to the decay\(^8\)

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}}V_{tb}V_{td}^* \left( \sum_{i=3}^{10} C_i O_i \right). \quad (1)$$

For convenience, we list below the operators in $H_{\text{eff}}$ for $b \to d\bar{s}s$:

\[
\begin{align*}
O_3 &= \bar{d}_\alpha \gamma^\mu Lb_\alpha \cdot \bar{s}_\beta \gamma^\mu Ls_\beta, & O_4 &= \bar{d}_\alpha \gamma^\mu Lb_\beta \cdot \bar{s}_\beta \gamma^\mu Ls_\alpha, \\
O_5 &= \bar{d}_\alpha \gamma^\mu Lb_\alpha \cdot \bar{s}_\beta \gamma^\mu Rs_\beta, & O_6 &= \bar{d}_\alpha \gamma^\mu Lb_\beta \cdot \bar{s}_\beta \gamma^\mu Rs_\alpha, \\
O_7 &= \frac{3}{2} \bar{d}_\alpha \gamma^\mu Lb_\alpha \cdot e_\gamma \bar{s}_\beta \gamma^\mu Rs_\beta, & O_8 &= \frac{3}{2} \bar{d}_\alpha \gamma^\mu Lb_\beta \cdot e_\gamma \bar{s}_\beta \gamma^\mu Rs_\alpha, \\
O_9 &= \frac{3}{2} \bar{d}_\alpha \gamma^\mu Lb_\alpha \cdot e_\gamma \bar{s}_\beta \gamma^\mu Ls_\beta, & O_{10} &= \frac{3}{2} \bar{d}_\alpha \gamma^\mu Lb_\beta \cdot e_\gamma \bar{s}_\beta \gamma^\mu Ls_\alpha.
\end{align*}
\]
Where $\alpha$ and $\beta$ are the $SU(3)$ color indices and $L = (1 - \gamma_5)/2$, $R = (1 + \gamma_5)/2$. The Wilson coefficients evaluated at $\mu = m_b$ scale are

$$
C_3 = 0.014, \quad C_4 = -0.035, \quad C_5 = 0.009, \quad C_6 = -0.041,
C_7 = -0.002/137, \quad C_8 = 0.054/137, \quad C_9 = -1.292/137, \quad C_{10} = 0.263/137,
$$

(3)

Using the effective Hamiltonian and naive factorization hypothesis, it is easy to write down the amplitude for $\bar{B}_d^0 \to \phi\gamma$

$$A(\bar{B}_d^0 \to \phi\gamma) = -\frac{G_F}{\sqrt{2}}V_{tb}V_{td}^* \left[ \left( C_3 + \frac{C_4}{N_c} + C_5 + \frac{C_6}{N_c} \right) - \frac{1}{2} \left( C_7 + \frac{C_8}{N_c} + C_9 + \frac{C_{10}}{N_c} \right) \right]
\times \sqrt{4\pi\alpha_e f_\phi M}\phi F_V \left\{ -\epsilon_{\mu\nu\rho\sigma} \eta_{l+}^\mu \epsilon_{l+}^\nu v^\rho q^\sigma + i \left( (\eta_{l+}^\mu \cdot \epsilon_{l+}^\nu)(v \cdot q) - (\eta_{l+}^\nu \cdot q)(v \cdot \epsilon_{l+}^\mu) \right) \right\}
$$

(4)

where $\eta_{l+}^\mu$ and $\epsilon_{l+}^\nu$ are transverse polarization vectors of $\phi$ and photon respectively. The form factor $F_V$ is defined by

$$
\langle \gamma(\epsilon^*, q)|\bar{d}\gamma_\mu(1 - \gamma_5)b|\bar{B}_d^0 \rangle = \sqrt{4\pi\alpha_e} \left[ -F_V \epsilon_{\mu\nu\rho\sigma} \epsilon^* \nu v^\rho q^\sigma + i F_A (\epsilon_\mu^* q \cdot v - q_\mu v \cdot \epsilon^*) \right].
$$

(5)

To leading power of $O(1/m_b)$, $F_V$ and $F_A$ read

$$
F_V = F_A = \frac{Q_d f_B M_B}{2\sqrt{2}E_\gamma} \int dl_+ \Phi_{B_1}(l_+)/l_+.
$$

(6)

The contributions of strong penguin operators arising from the renormalization group evolution from the scale $\mu = M_W$ to $\mu = m_b$ is very small due to the cancellations between them: $C_3 \simeq -C_4/3$ and $C_5 \simeq -C_6/3$. Obviously the amplitude is dominated by electro-weak penguin(EWP).

Now we can write down the helicity amplitude

$$
\mathcal{M}_{++,} = i \frac{G_F}{\sqrt{2}} V_{tb} V_{td}^* \sqrt{4\pi\alpha_e} F_V f_\phi m_\phi M_B \left[ a_3 + a_5 - \frac{1}{2} (a_7 + a_9) \right],
$$

(7)

$$
\mathcal{M}_{--} = 0,
$$

(8)

where $a_i = C_i + C_{i+1}/N_C$. It is interesting to note that the $\phi$ meson and the photon in the decay are right-handed polarized in the SM. It is also easy to realize that the decay is very rare because of helicity suppression as well as small $V_{td}$, $f_B$ and $f_\phi$. Using $f_\phi = 254 MeV$, $f_B = 180 MeV$, $|V_{td}| = 0.008$, $N_C = 3$ and $\lambda_B^{-1} = \int dl_+ \Phi_{B_1}(l_+)/l_+ = (0.35 GeV)^{-1} \frac{3}{4}$, we get

$$
\mathcal{B}(\bar{B}_d^0 \to \phi\gamma) = 3.5 \times 10^{-13}.
$$

(9)

4
In the above calculations, non-factorizable contributions are neglected. However, the leading non-factorizable diagrams in Fig.2 should be taken into account. For this purpose, the QCD factorization framework[6] invented recently by Beneke, Buchalla, Neubert and Sachrajda is very suitable. The framework incorporate many important theoretical aspects of QCD like color transparency, heavy quark limit and hard-scattering which allows us to calculate nonfactorizable contributions systematically.

To calculate the nonfactorizable diagrams as depicted by Fig.2, we take photon and $\phi$ meson flying along $n_− = (1, 0, 0, −1)$ and $n_+ = (1, 0, 0, 1)$ directions respectively. We need the two-particle light-cone projector for B meson and $\phi$ meson

$$M^B_{\alpha\beta} = \frac{i}{4N_c} f_B M_B \left\{ (1 + \not{v}) \gamma_5 \left[ \Phi_{B1}(l_+) + \not{\not{n}}_\perp \Phi_{B2}(l_+) \right] \right\}_{\alpha\beta},$$

(10)

$$M^\phi_{\delta\gamma} = -\frac{f_\phi m_\phi}{4N_c} \left\{ \gamma_\perp \gamma_\perp \left( u + i \epsilon_{\mu\rho\sigma} \epsilon_\perp^{\nu} n_\perp^{\mu} n_\perp^{\sigma} \gamma_5 \frac{g_\perp^{(a)}(u)}{8} \right) \right\} \delta_{\delta\gamma},$$

(11)

which encode the relevant non-perturbative bound state dynamics of the initial B meson and the final $\phi$ meson. $\Phi_{B1}(l_+)$ and $\Phi_{B2}(l_+)$ are the leading twist light-cone distribution function of B meson[10]. $g_\perp^{(a)}(u)$ and $g_\perp^{(a)}(u)$ are twist-3 distribution amplitude of $\phi$ meson[9], $g_\perp^{(a)}(u) = dg_\perp^{(a)}(u)/du$. The detail discussions on these projectors could be found in Refs.[9, 10, 11].

The amplitude for $B^0_d \to \phi\gamma$ can be presented by the schematic formula

$$\langle \phi\gamma | H_{\text{eff}} | B \rangle = -\frac{G_F}{\sqrt{2}} V_{tb} V_{td}^\ast \left[ (a_3' + a_5') - \frac{1}{2} (a_7' + a_9') \right]$$

$$\times \sqrt{4\pi\alpha_e f_\phi m_\phi F_V} \left\{ -\epsilon_{\mu\rho\sigma} \not{n}_\perp^{\mu} \epsilon_\perp^{\nu} v^\rho q^\sigma + i \left[ (\not{n}_\perp^{\mu} \cdot v \cdot q) - (\not{n}_\perp^{\mu} \cdot q) (v \cdot \epsilon_\perp^{\nu}) \right] \right\}$$

(12)

The $O(\alpha_s)$ corrections are summarized in $a_i'$s which are calculated to be

$$a_3' = a_3 + \frac{\alpha_s}{4\pi N_C f_\phi} C_F C_4 F_1,$$

(13)

$$a_5' = a_5 + \frac{\alpha_s}{4\pi N_C f_\phi} C_F C_6 F_2,$$

(14)

$$a_7' = a_7 + \frac{\alpha_s}{4\pi N_C f_\phi} C_F C_8 F_2,$$

(15)

$$a_9' = a_9 + \frac{\alpha_s}{4\pi N_C f_\phi} C_F C_{10} F_1,$$

(16)

where $F_{1,2}$ arise from one gluon exchange between the two currents of color-octet penguin
Figure 2: Nonfactorizable diagrams for $\bar{B}_d^0 \to \phi \gamma$. Other diagrams with the photon radiating from the heavy b quark and the energetic quark lines are suppressed.

operator $O_{4,6,8,10}$ as shown by fig.2.

\[
F_1 = \int_0^1 du \left( \frac{g_1^{(a\gamma)}(u)}{4} - g_1^{(\gamma)}(u) \right) \left[ -14 - 3i\pi - 12 \ln \frac{\mu}{m_b} 
+ \left( 5 + \frac{u}{1-u} \ln u \right) - \frac{\pi^2}{3} + 2Li_2\left( \frac{u-1}{u} \right) \right],
\]

\[
F_2 = \int_0^1 du \left( g_1^{(\gamma)}(u) + \frac{g_1^{(a\gamma)}(u)}{4} \right) \left[ -14 - 3i\pi - 12 \ln \frac{\mu}{m_b} 
+ \left( 5 + \frac{u}{1-u} \ln u \right) - \frac{\pi^2}{3} + 2Li_2\left( \frac{u-1}{u} \right) \right].
\]

In the calculation, the $\overline{MS}$ renormalization scheme is used. We have neglected the small effect of box diagrams and the diagrams with photon radiating from energetic strange quarks, which are further suppressed by $\Lambda_{QCD}/M_B$. We also have neglect $l_+^2$ terms entered in the loop calculation which are the higher twist effect, in this way, the integral involved $\Phi_{B^0}(l_+)$ absents and the remain integrals are related to the form factor $F_V$.

Including $O(\alpha_s)$ contributions, the helicity amplitude is

\[
\mathcal{M}_{++} = i \frac{G_F}{\sqrt{2}} V_{tb} V_{td}^* \sqrt{4\pi\alpha_e} F_V f_{\phi} m_{\phi} M_B \left[ a'_3 + a'_5 - \frac{1}{2}(a'_7 + a'_9) \right],
\]
\[ M_- = 0. \] (19)

To give numerical result, we take \( \mu = m_b \) and \( f_\phi^+ = 215 \text{MeV} \). The branching ratio is estimated to be

\[ B(\bar{B}_d^0 \rightarrow \phi\gamma) = 3.6 \times 10^{-12}. \] (20)

It indicates that the decay is too rare to be measured at the running B factories BaBar and BELLE. However, the decay may be accessible at LHCb and the planning super high luminosity B factories at KEK and SLAC. Furthermore, it could be enhanced by New Physics, and the enhancement might be large enough for being measured at these facilities.

3 \( \bar{B}_d^0 \rightarrow \phi\gamma \) as a probe of R parity violation

In the minimal supersymmetric standard model (MSSM) a discrete symmetry called R-parity is invoked to forbid gauge invariant lepton and baryon number violating operators. The R-parity of a particle is given by

\[ R_p = (-1)^{L+2S+3B}, \]

where \( L \) and \( B \) are lepton and baryon numbers, and \( S \) is the spin. However there is no deep theoretical motivation for imposing R-parity and it is interesting to explore the phenomenology of R-parity violation.

We start our exploration from the R-parity violating superpotential

\[ W_R = \frac{1}{2} \lambda LLE^c + \lambda' LQD^c + \frac{1}{2} \lambda'' U^c D^c D^c. \] (21)

We are interested in the \( \lambda' \) and \( \lambda'' \) terms since they are relevant to the process \( b \rightarrow d s \bar{s} \). Writing down all indices explicitly, it reads

\[ W_R = \varepsilon^{ab} \delta^{\alpha\beta} \lambda'_{ijk} L_{ia} Q_{jb} D^c_{k\beta} + \frac{1}{2} \varepsilon^{\alpha\beta\gamma} \left[ \lambda''_{i[jk]} U^c_{i\alpha} D^c_{j\beta} D^c_{k\gamma} \right], \] (22)

where \( U^c_{i\alpha} \), \( D^c_{j\beta} \) and \( D^c_{k\gamma} \) are the superfields of the right handed quarks and/or squarks respectively, the superscript \( c \) denotes charge conjugation. \( i, j \) and \( k \) are generation indices and \( \alpha, \beta \) and \( \gamma \) are SU(3) color triplet indices. It follows from the antisymmetry of \( \varepsilon^{\alpha\beta\gamma} \) that \( \lambda''_{i[jk]} \) is antisymmetric in the last two indices.

In terms of four components Dirac spinor, from \( W_R \) we can read

\[ \mathcal{L}^R_{eff} = \frac{1}{2} \varepsilon^{\alpha\beta\gamma} \left[ \lambda''_{i[jk]} \bar{u}_{Ri\alpha} \left( \bar{d}_{j\beta} R d_{k\gamma} - \{ j \leftrightarrow k \} \right) \right] + \lambda'_{i[jk]} \bar{\nu}_{Li\alpha} \bar{d}_{k\gamma} L d_{j} + h.c. \] (23)
From $\mathcal{L}_{\text{eff}}$, we get the effective Hamiltonian for $b \to d s \bar{s}$

$$
\mathcal{H}_{\text{R}} = \frac{1}{m_{u_i}^2} \varepsilon^{\alpha\beta\gamma} \varepsilon^{\alpha'\beta'\gamma'} \left[ \lambda''_{123} \lambda''_{112} (s_\gamma R S) (s\gamma L d) + \lambda''_{123} \lambda_{121} (s_\gamma R S) (d\gamma L s) 
+ \lambda''_{231} \lambda_{212} (b_\beta R S) (s\gamma L d) + \lambda''_{132} \lambda_{121} (b_\beta R S) (d\gamma L s) 
+ \lambda''_{132} \lambda_{121} (a_\alpha R S) (s \beta L b) + \lambda''_{123} \lambda_{121} (a_\alpha R S) (s \beta L b) \right].
$$

(24)

Contracting $\varepsilon^{\alpha\beta\gamma} \varepsilon^{\alpha'\beta'\gamma'}$ and performing Fierz transformations, we get

$$
\mathcal{H}_{\text{R}}(b \to d s \bar{s}) = - \frac{2}{m_{u_i}^2} \lambda''_{231} \lambda''_{112} \left[ (s_\gamma R S)(d\gamma' L S) - (s_\gamma R S)(d\gamma' L S) \right] - \frac{1}{2m_{u_i}^2} \lambda''_{132} \lambda_{121} (s_\gamma R S) (d\gamma' L S) + \lambda''_{132} \lambda_{121} (s_\gamma R S) (d\gamma' L S)
+ \lambda''_{132} \lambda_{121} (d_\alpha R S) (s \gamma' L S) + \lambda''_{123} \lambda_{121} (d_\alpha R S) (s \gamma' L S)
$$

(25)

where we have used the relation

$$
(q_\gamma R L q^c) = -(q_\gamma R L q).
$$

(26)

The effective Hamiltonian is near to the idea form. Using renormalization group to run it from sfermion mass scale $m_{\tilde{f}}$ (100GeV assumed) down to $m_b$ scale, it reads

$$
\mathcal{H}_{\text{R}}(b \to d s \bar{s}) = - \frac{2}{m_{u_i}^2} \eta^{-4/\beta_0} \lambda''_{231} \lambda''_{112} \left[ (s_\gamma R S)(d\gamma' L S) - (s_\gamma R S)(d\gamma' L S) \right] - \frac{1}{2m_{u_i}^2} \eta^{-8/\beta_0} \lambda''_{132} \lambda_{121} (s_\gamma R S) (d\gamma' L S) + \lambda''_{132} \lambda_{121} (s_\gamma R S) (d\gamma' L S)
+ \lambda''_{132} \lambda_{121} (d_\alpha R S) (s \gamma' L S) + \lambda''_{123} \lambda_{121} (d_\alpha R S) (s \gamma' L S)
$$

(27)

with $\eta = \frac{\alpha_s(m_{\tilde{f}})}{\alpha_s(m_b)}$ and $\beta_0 = 11 - \frac{2}{3} n_f$.

Now we are ready to write down R-parity violation contributions in $\tilde{B}_d^0 \to \phi \gamma$ decays

$$
\mathcal{M}^R(\tilde{B}_d^0 \to \phi \gamma) = \sqrt{4\pi \alpha_s} F_V f_{\phi} m_{\phi} \left[ i(\eta^* \cdot \epsilon^*)(u \cdot q) + \epsilon_{\mu \nu \alpha \beta} \eta^* \epsilon^* \epsilon \cdot q \right]
\times \left\{ \frac{1}{2m_{\tilde{u}}^2} \alpha_s \lambda''_{231} \lambda''_{112} \eta^{-4/\beta_0} \left( 1 - \frac{1}{N_C} - \frac{\alpha_s}{4\pi} \frac{C_F}{N_C} F_1 \right) + \frac{1}{8m_{\tilde{u}}^2} \alpha_s \lambda''_{121} \lambda''_{121} \eta^{-8/\beta_0} \left( \frac{1}{N_C} + \frac{\alpha_s}{4\pi} \frac{C_F}{N_C} F_2 \right) \right\}
- \sqrt{4\pi \alpha_s} F_V f_{\phi} m_{\phi} \left[ i(\eta^* \cdot \epsilon^*)(u \cdot q) - \epsilon_{\mu \nu \alpha \beta} \eta^* \epsilon^* \epsilon \cdot q \right]
\times \frac{1}{8m_{\tilde{u}}^2} \alpha_s \lambda''_{132} \lambda''_{112} \eta^{-8/\beta_0} \left( \frac{1}{N_C} + \frac{\alpha_s}{4\pi} \frac{C_F}{N_C} F_1 \right).
$$

(28)
We note that only $u$ channel squarks and sneutrino mediated terms contribute because of $\langle \gamma | \bar{d} (1 \pm \gamma_5) b | \bar{B}^0_d \rangle = 0$. Immediately, we derive

$$
\mathcal{M}^R_{++} = i M_B \sqrt{4 \pi \alpha_e F_V f_\phi m_\phi} \frac{1}{8m_{\tilde{\nu}_i}^2} \lambda'_{i32} \lambda'^*_{i12} \gamma^{-8/\beta_0} \left( \frac{1}{N_C} + \frac{\alpha_s C_F}{4\pi N_c} F_1 \right),
$$

(29)

$$
\mathcal{M}^R_{--} = -i M_B \sqrt{4 \pi \alpha_e F_V f_\phi m_\phi} \left\{ \frac{1}{2m_{\tilde{u}_i}} \lambda''_{i23} \lambda''^*_{i12} \gamma^{-4/\beta_0} \left( 1 - \frac{1}{N_C} - \frac{\alpha_s C_F}{4\pi N_c} F_1 \right) + \frac{1}{8m_{\tilde{\nu}_i}^2} \lambda'_{i21} \lambda'^*_{i23} \gamma^{-8/\beta_0} \left( \frac{1}{N_C} + \frac{\alpha_s C_F}{4\pi N_c} F_2 \right) \right\}.
$$

(30)

Figure 3: The branching ratio of $\bar{B}^0_d \to \phi \gamma$ as a function of R-parity violation couplings $|\lambda''_{i23} \lambda''_{i12}|$ (the thick solid curve), $|\lambda'_{i32} \lambda'^*_{i12}|$ (the thin solid curve) and $|\lambda'_{i21} \lambda'^*_{i23}|$ (the dash curve) respectively. The horizontal line is the SM prediction.

To present numerical results, we assume that only one sfermion contribute one time and universal mass 100GeV for sfermions $\tilde{u}_{Ri}$. Our results for R parity violation contributions are summarized in fig.3. We note that the decay is very sensitive to R parity violating couplings. Within the available lowest upper bounds for $|\lambda''_{i23} \lambda''_{i12}|$ [15, 16, 17],

$$
|\lambda''_{i23} \lambda''_{i12}| < 6 \times 10^{-5} \left( \frac{m_{\tilde{u}_{Ri}}}{100} \right)^2, \quad |\lambda'_{i32} \lambda'^*_{i12}| < 4 \times 10^{-4} \left( \frac{m_{\tilde{\nu}_Li}}{100} \right)^2,
$$

(31)

$\mathcal{B}(\bar{B}^0_d \to \phi \gamma)$ could be enhanced to $10^{-9} \sim 10^{-8}$ which may be inaccessible at BELLE and BaBar. However, it is large enough for LHCb, BTeV and the planning super high luminosity B factories at KEK and SLAC.
Conclusion

We have studied pure penguin radiative annihilation process $\bar{B}_d^0 \to \phi \gamma$ by using QCD factorization for the hadronic dynamics. We find that non-factorizable contributions are larger than factorizable contributions in $\bar{B}_d^0 \to \phi \gamma$ decays. We estimate $B(\bar{B}_d^0 \to \phi \gamma) = 3.6 \times 10^{-12}$ in the SM. The smallness of these decays in the SM makes it a sensitive probe of flavor physics beyond the SM. To explore the potential, we have estimated the contribution of RPV couplings to this decay. We have found that the decay $\bar{B}_d^0 \to \phi \gamma$ is very sensitive to the $\lambda'' U^c D^c D^c$ and $\lambda' L^c D^c$ terms in the R-parity violating superpotential. Within the upper limits for $\lambda''_{23} \lambda''_{121}$ and $\lambda'_3 \lambda''_{12}$, it is found that $B(\bar{B}_d^0 \to \phi \gamma)$ could be enhanced to order of $10^{-9} \sim 10^{-8}$. We also note that the $\phi$ meson and the photon are right-handed polarized in the SM, but they can be left-handed polarized in RPV supersymmetry. We find that $\lambda'_3 \lambda''_{12}$ RPV couplings contribute to the right-handed polarized magnitude while $\lambda''_{23} \lambda''_{121}$ and $\lambda'_{121} \lambda'_2$ RPV couplings responsible to the left-handed polarized magnitude. Theoretically, measuring the polarization of photon (or the $\phi$ meson) can be used to probe the relative strength of RPV couplings. However, the interesting phenomena is faded by the small branching ratio of the exotic decay. In literature, there are interesting proposals for probing new physics by measuring the photon polarization in radiative decays $B \to K^* \gamma$\[^{18}\] . It always needs large amount of experimental data. For the more exotic decay $\bar{B}_d^0 \to \phi \gamma$, this phenomena will be remained academic unless the continuum background could be suppressed efficiently in data analyzing at the future B facilities. The interesting decay may be too rare to be accessible at Babar and Belle. However, the decay rate might be studied by LHCb at CERN, BTeV at Fermilab and the planning super high luminosity B factories at KEK and SLAC to probe new physics.

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