Probing new physics in $B \to f_0(980)K$ decays

A. K. Giri$^1$, B. Mawlong$^2$ and R. Mohanta$^2$

$^1$ Department of Physics, Punjabi University, Patiala - 147002, India
$^2$ School of Physics, University of Hyderabad, Hyderabad - 500 046, India

Abstract

We study the hadronic decay modes $B^{\pm(0)} \to f_0(980)K^{\pm(0)}$, involving a scalar and a pseudoscalar meson in the final state. These decay modes are dominated by the loop induced $b \to s\bar{q}q$ ($q = s, u, d$) penguins along with a small $b \to u$ tree level transition (for $B^+ \to f_0K^+$) and annihilation diagrams. Therefore, the standard model expectation of direct CP violation is negligibly small and the mixing induced CP violation parameter in the mode $B^0 \to f_0K_S$ is expected to give the same value of $\sin(2\beta)$, as extracted from $B^0 \to J/\psi K_S$ but with opposite sign. Using the generalized factorization approach we find the direct CP violation in the decay mode $B^+ \to f_0K^+$ to be of the order of few percent. We then study the effect of the R-parity violating supersymmetric model and show that the direct CP violating asymmetry in $B^+ \to f_0(980)K^+$ could be as large as $\sim 80\%$ and the mixing induced CP asymmetry in $B^0 \to f_0K_S$ (i.e., $-S_{f_0K_S}$) could deviate significantly from that of $\sin(2\beta)_{J/\psi K_S}$.

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I. INTRODUCTION

The currently running $B$ factories, such as Belle and Babar, are providing us huge data in the b-quark sector. The main objectives of these $B$ factories is to critically test the standard model (SM) predictions and to look for possible signature of new physics. For this purpose, a variety of useful observables are being measured and are compared with the corresponding theoretical predictions. One of the important observables of this kind is the CP asymmetry parameter in various $B$ meson decays and the other one being the branching ratio for rare $B$ processes.

Recently, both Belle\cite{1,2,3,4} and Babar\cite{5,6,7,8,9} have reported the measurement of branching ratios and CP violating parameters in the rare decay modes $B^0, + \rightarrow f_0(980)K^0, +$, involving a scalar and a pseudoscalar meson in the final state. The measured decay rates for the mode $B^+ \rightarrow f_0K^+$ are

$$ BR(B^+ \rightarrow f_0(980)K^+ \rightarrow \pi^+\pi^-K^+) = (8.78 \pm 0.82^{+0.85}_{-1.76}) \times 10^{-6}, \quad [3] $$

$$ BR(B^+ \rightarrow f_0(980)K^+ \rightarrow \pi^+\pi^-K^+) = (9.47 \pm 0.97^{+0.62}_{-0.88}) \times 10^{-6}, \quad [8] \quad (1) $$

with an average

$$ BR(B^+ \rightarrow f_0(980)K^+ \rightarrow \pi^+\pi^-K^+) = (9.21 \pm 0.97) \times 10^{-6}. \quad (2) $$

For the process $B^0 \rightarrow f_0K^0$, the measured rates are

$$ BR(B^0 \rightarrow f_0(980)K^0 \rightarrow \pi^+\pi^-K^0) = (7.60 \pm 1.66^{+0.76}_{-0.89}) \times 10^{-6}, \quad [4] $$

$$ BR(B^0 \rightarrow f_0(980)K^0 \rightarrow \pi^+\pi^-K^0) = (5.5 \pm 0.7 \pm 0.7 \times 10^{-6} \quad [9] \quad (3) $$

The absolute branching ratios for $B \rightarrow f_0K$ processes depend on the branching fraction of $f_0 \rightarrow \pi^+\pi^-$ process. Using the results from\cite{16} for $\Gamma(f_0 \rightarrow \pi\pi) = 64 \pm 8$ MeV, $\Gamma_{f_0}^{tot} = 80 \pm 10$ MeV along with the relation $\Gamma(f_0 \rightarrow \pi^+\pi^-) = \frac{2}{3}\Gamma(f_0 \rightarrow \pi\pi)$, we obtain the branching ratios for $B \rightarrow f_0K$ processes as

$$ BR(B^+ \rightarrow f_0(980)K^+) = (17.38 \pm 3.47) \times 10^{-6}, $$

$$ BR(B^0 \rightarrow f_0(980)K^0) = (11.26 \pm 2.52) \times 10^{-6}. \quad (4) $$

The mixing induced parameter for the process $B^0 \rightarrow f_0K_S$, observed by both Babar and
Belle as
\[ \sin(2\beta)_{f_0K_S} = 0.95^{+0.23}_{-0.32} \pm 0.10, \quad [10] \]
\[ \sin(2\beta)_{f_0K_S} = 0.18 \pm 0.23 \pm 0.11, \quad [11] \]
(5)

with an average
\[ \sin(2\beta)_{f_0K_S} = 0.51 \pm 0.19, \quad (6) \]
which has nearly one sigma deviation from that of \( \sin(2\beta)_{b \to c\bar{s}s} = 0.687 \pm 0.032 \) \[12\]. These observations not only provide us another way to test the SM and/or to look for new physics but also may help us to understand the nature of the light scalar meson \( f_0(980) \). It should be noted here that the mixing induced CP violation parameter seems to be, at present, not deviated significantly from its SM expectation. But, since the error bars are quite large the situation is still very much conducive to explore some non-standard physics.

The light scalar mesons with masses below 1 GeV is considered as a controversial issue for a long time. Even today, there exists no consensus on the nature of the \( f_0(980) \) and \( a_0(980) \) mesons. While the low-energy hadron phenomenology has been successfully understood in terms of the constituent quark model, the scalar mesons are still puzzling and the quark composition of the light scalar mesons are not understood with certainty. The structure of the scalar meson \( f_0(980) \) has been discussed for decades and appears to be still not clear. There were attempts to interpret it as \( K\bar{K} \) molecular states \[13\], four quark states \[14\] and normal \( q\bar{q} \) states \[15\]. However, recent studies of \( \phi \to \gamma f_0 \) \( (f_0 \to \gamma\gamma) \) \[16, 17\] and \( D_s^+ \to f_0\pi^+ \) decays \[18\] favor the \( q\bar{q} \) model. Since \( f_0(980) \) is produced copiously in \( D_s \) decays, this supports the picture of large \( s\bar{s} \) component in its wave function, as the dominant mechanism in the \( D_s \) decay is \( c \to s \) transition. The prominent \( s\bar{s} \) nature of \( f_0(980) \) has been supported by the radiative decay \( \phi \to f_0(980)\gamma \) \[19\]. In this interpretation, the flavor content of \( f_0 \) is given by \( f_0 = n\bar{n}\sin\theta + s\bar{s}\cos\theta \) with \( n\bar{n} = (u\bar{u} + d\bar{d})/\sqrt{2} \). A mixing angle of \( \theta = 138^\circ \pm 6^\circ \) has been experimentally determined from \( \phi \to \gamma f_0 \) decays \[16\]. We will follow this structure for our study in this paper.

Theoretically, these decay modes have been studied in the standard model using perturbative QCD \[20\] and QCD factorization approach \[21, 22\]. In this paper, we would like to study the decay modes \( B^0 \to f_0(980)K^0 \) and \( B^+ \to f_0(980)K^+ \) using the generalized factorization approach. We consider \( f_0(980) \) to be composed of \( f_0(980) = n\bar{n}\sin\theta + s\bar{s}\cos\theta \) with
dominant $s\bar{s}$ composition. Therefore, these processes may be considered, at the leading order, as dominated by $b \rightarrow s\bar{s}s$ penguin amplitudes. Hence, the mixing induced CP violation in the decay mode $B^0 \rightarrow f_0 K_S$ is expected to give the same value of $\sin(2\beta)$ as extracted from $B^0 \rightarrow J/\psi K_S$, with an uncertainty of 5%. Comparison of these two values, therefore, could be a sensitive probe for physics beyond the SM. Since the predicted branching ratios available from previous studies [20, 21, 22] are not in agreement with the experimental values, we would like to see the effect of R-parity violating supersymmetric (RPV) model in these modes. Moreover, since in this paper we are interested to see whether it is possible to extract any signature of new physics (NP) from these modes or not, we resort ourselves to generalized factorization approach in analyzing these modes.

The paper is organized as follows. In section II, we analyze these modes in the standard model. The basic formula for CP violating parameters are presented in section III. The contributions arising from R-parity violating model are presented in section IV and section V contains our conclusion.

II. STANDARD MODEL CONTRIBUTION

The effective Hamiltonian describing the charmless hadronic $B$ decays is given as

$$H_{eff} = \frac{G_F}{\sqrt{2}} \left[ V_{ub}^* V_{us} \sum_{i=1}^{2} C_i O_i - V_{tb}^* V_{ts} \sum_{j=3}^{10} C_j O_j \right].$$

where $G_F$ is the Fermi coupling constant, $C_i$'s are the Wilson coefficients, $O_{1,2}$ are the tree operators and $O_{3-10}$ are QCD and electroweak penguin operators.

To calculate the branching ratios of the $B \rightarrow f_0 K$ decay processes, we adopt the generalized factorization framework to evaluate the hadronic matrix elements i.e., $\langle O_i \rangle = \langle f_0 K | O_i | B \rangle$. In this approximation, these hadronic matrix elements can be parametrized in terms of the decay constants and the form factors which are defined as

$$\langle 0 | A^\mu | K(k) \rangle = i f_K k^\mu, \quad \langle 0 | \bar{q}q | f_0 \rangle = m_{f_0} f_{f_0}^q,$$

$$\langle K(k) | (V-A)_\mu | B(P) \rangle = \left[ (P+k)_\mu - \left( \frac{m_B^2 - m_K^2}{q^2} \right) q_\mu \right] F^BK_1(q^2)$$

$$+ \left( \frac{m_B^2 - m_K^2}{q^2} \right) q_\mu F^BK_0(q^2),$$

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\[ \langle f_0(q) | (V - A)_{\mu} | B(P) \rangle = i \left\{ \left[ (P + q)_{\mu} - \left( \frac{m_B^2 - m_{f_0}^2}{k^2} \right) k_{\mu} \right] F_1^{Bf_0}(k^2) \right. \\
+ \left. \left( \frac{m_B^2 - m_{f_0}^2}{k^2} \right) k_{\mu} F_0^{Bf_0}(k^2) \right\}, \]  

where \( V \) and \( A \) denote the vector and axial-vector currents, \( f_K \) and \( f_{f_0} \) are the decay constants of \( K \) and \( f_0 \) mesons, \( F_{0,1}(q^2) \) are the form factors and \( P, q, k \) are the momenta of \( B, f_0 \) and \( K \) mesons satisfying the relation \( q = P - k \).

Now let us first consider the process \( B^+ \rightarrow f_0 K^+ \). Within the SM it receives contribution from \( b \rightarrow u \) tree, \( b \rightarrow s \bar{q} q \) (with \( q = u, s \)) penguins and annihilation diagrams. Using Eqs. (7)-(10), one can obtain the amplitude in the SM as

\[ A(B^+ \rightarrow f_0 K^+) = -\frac{G_F}{\sqrt{2}} \left\{ \left[ V_{ub}^{*}V_{us}a_1 - V_{tb}^{*}V_{ts}(a_4 + a_{10} - r_\chi (a_6 + a_8)) \right] X \\
- V_{tb}^{*}V_{ts}(2a_6 - a_8)Y - \left[ V_{ub}^{*}V_{us}a_1 \right. \\
- V_{tb}^{*}V_{ts}(a_4 + a_{10} - \left. \frac{2(a_6 + a_8)m_B^2}{(m_b + m_u)(m_s + m_u)} \right] \right\}, \]  

(11)

where

\[ r_\chi = \frac{2m_K^2}{(m_b + m_u)(m_s + m_u)}, \quad X = f_K(m_B^2 - m_{f_0}^2)F_0^{Bf_0}(m_K^2), \]

\[ Y = f_{f_0}^{*} m_{f_0} \frac{m_B^2 - m_K^2}{m_b - m_s} F_0^{B_K}(m_{f_0}^2), \quad Z = f_B(m_{f_0}^2 - m_K^2)F_0^{f_0K}(m_B^2), \]

(12)

and \( a_i ' s \) are the combinations of Wilson coefficients given by

\[ a_{2i-1} = C_{2i-1} + \frac{1}{N_C} C_{2i}, \quad a_{2i} = C_{2i} + \frac{1}{N_C} C_{2i-1}, \quad (i = 1, 2, 3, 4, 5) \]

(13)

with \( N_C \) as the number of colors.

The corresponding neutral process \( B^0 \rightarrow f_0 K^0 \) receives contribution only from \( b \rightarrow s \bar{q} q \) (with \( q = s, d \)) penguins and annihilation diagrams. Thus, one can write the amplitude\(^1\) for this process as

\[ A(B^0 \rightarrow f_0 K^0) = \frac{G_F}{\sqrt{2}} V_{tb}^{*}V_{ts} \left\{ \left[ a_4 - \frac{a_{10}}{2} - r_\chi (a_6 - \frac{a_8}{2}) \right] X \\
+ (2a_6 - a_8)Y - \left[ a_4 - \frac{a_{10}}{2} - \frac{(2a_6 - a_8)m_B^2}{(m_b + m_d)(m_s + m_d)} \right] Z \right\}, \]

(14)

\(^1\) The sign of the coefficients of \( Y \) in Eqs. (11) and (14) are found to be opposite to that of Ref. [28].
where $r_{\chi_1}$ can be obtained from $r_{\chi}$ by replacing the $K^+$ and $u$-quark masses by $K^0$ and $d$ masses. The branching ratios can be obtained from these amplitudes as

$$BR(B \to f_0 K) = \frac{|p_c| \tau_B}{8\pi m_B^2} |A(B \to f_0 K)|^2,$$

where $|p_c|$ is the c.m. momentum of the final mesons and $\tau_B$ is the lifetime of the $B$ meson. For numerical analysis, we use the particle masses and lifetimes from [23]. The current quark masses are taken as $m_b = 4.88$ GeV, $m_s = 122$ MeV, $m_d = 7.6$ MeV and $m_u = 4.2$ MeV. The values of the effective QCD parameters ($a_i$'s) are taken from [24], which are evaluated at the scale $\mu = m_b/2$. For the CKM matrix elements, we use the Wolfenstein parametrization with the parameters $A=0.801$, $\lambda = 0.2265$, $\bar{\rho} = 0.189$ and $\bar{\eta} = 0.358$ [25]. The form factors describing the transition $B \to f_0$ are given as [21]

$$F_{B^-f_0} = \frac{1}{\sqrt{2}} \sin \theta F_{B^-f_0}^{u\bar{u}}, \quad F_{B^0f_0}^{f_0} = \frac{1}{\sqrt{2}} \sin \theta F_{B^0f_0}^{d\bar{d}},$$

with $F_{B^0f_0}^{f_0}(0) \ (q\bar{q} = u\bar{u}$ or $d\bar{d})$ being of the order of 0.25 [26]. For the $q^2$ dependence, we assume the simple pole dominance as

$$F_{0}^{Bf_0}(q^2) = \frac{F_{0}^{Bf_0}(0)}{1 - q^2/m_P^2},$$

with $m_P$ being the mass of the $0^-$ pole state with the same quark content as the current under consideration. For the form factors, describing $B \to K$ transition, we use the corresponding QCD sum rule value [27]

$$F_{0}^{BK}(m_{f_0}^2) = \frac{0.3302}{1 - \frac{m_{f_0}^2}{3.46}}.$$ (18)

The annihilation form factor $F_{0}^{f_0K}(q^2)$ is expected to be suppressed at large momentum transfer (i.e., $q^2 = m_B^2$) due to helicity suppression. However, it may receive long distance contributions from nearby resonances via final state interactions. In Ref. [28], its value is extracted using the experimental values of $BR(B \to f_0 K)$, where it has been shown that in order to explain the observed data in the SM one requires large value of annihilation form factor, if the $B \to f_0$ form factor will be $F_{0}^{Bf_0} \leq 0.2$. Since, we are interested to look for new physics signature in this mode, here we use the lowest value of $|F_{0}^{f_0K}|$, which is around 0.03, as seen from figure-2 of [28]. Furthermore, since both the components of $f_0$ ($n\bar{n}$ and $s\bar{s}$) are involved in the annihilation topology, the corresponding amplitude should be multiplied by $(\sin \theta/\sqrt{2} + \cos \theta)$. 

\[ \text{6} \]
The decay constants used are $f_K=0.16$ GeV, $f_B=0.19$ GeV and $\tilde{f}_s = \frac{m^{(s)}_{f_0}}{m_{f_0}} \tilde{f}_s \cos \theta$ with $\tilde{f}_s(\mu = 2.1\text{GeV})=0.39$ GeV and $m^{(s)}_{f_0} \simeq (1.02 \pm .05)$ GeV [21].

Using these values and the mixing angle $\theta = 138^\circ$, we obtain the branching ratios for the $B \to f_0(980)K$ processes as

$$BR(B^+ \to f_0(980)K^+) = 6.56 \times 10^{-6},$$
$$BR(B^0 \to f_0(980)K^0) = 4.73 \times 10^{-6},$$

which are quite below the experimental values (4). The variation of the branching ratios for the strange, non-strange mixing angle $\theta$ between 0 and $\pi$ are shown in figures 1 and 2. Thus, one can see from the figure-1 that for $B^+ \to f_0 K^+$ process generalized factorization approach cannot accommodate the experimental data for any value of the mixing angle $\theta$. For the $B^0 \to f_0 K^0$ mode also it cannot explain the data unless $\theta$ is very close to 0 or $\pi$ as seen from figure-2.

![Figure 1: The branching ratio for the process $B^- \to f_0(980)K^-$ (in units of $10^{-6}$), versus the mixing angle $\theta$ in degrees.](image)

### III. CP VIOLATION PARAMETERS

Here, we briefly present the basic and well known formula for the CP violating parameters. Let us first consider the process $B^+ \to f_0 K^+$, which has only direct CP violation. The amplitude for this process can be symbolically written as

$$\mathcal{A}(B^+ \to f_0 K^+) = \lambda_u^* |A_u| e^{i\delta_u} + \lambda_t^* |A_t| e^{i\delta_t},$$
$$\mathcal{A}(B^0 \to f_0 K^0) = \lambda_u |A_u| e^{i\delta_u} + \lambda_t |A_t| e^{i\delta_t},$$

(20)
where $\lambda_q = V_{qb}V_{qs}^*$ with ($q = u, t$) denote the product of CKM matrix elements which contain
the weak phase information. It should be noted that the weak phase of $\lambda_u$ is $\arg(V_{ub}V_{us}) = \gamma$
and that of $\lambda_t$ is $\arg(V_{tb}V_{ts}) = \pi$. $A_u$ and $A_t$ denote the contributions arising from the
current operators proportional to $\lambda_u$ and $\lambda_t$ respectively and the corresponding strong phases
are taken as $\delta_u$ and $\delta_t$.

For the charged $B^\pm \rightarrow f_0K^\pm$ decays the CP violating rate asymmetry in the partial rates
is defined as follows:

$$A_{\text{CP}} = \frac{\Gamma(B^+ \rightarrow f_0K^+) - \Gamma(B^- \rightarrow f_0K^-)}{\Gamma(B^+ \rightarrow f_0K^+) + \Gamma(B^- \rightarrow f_0K^-)}$$

$$= \frac{2r \sin \gamma \sin(\delta_u - \delta_t)}{1 + r^2 - 2r \cos \gamma \cos(\delta_u - \delta_t)}, \quad (21)$$

where $r = |\lambda_u A_u/\lambda_t A_t|$. Thus to obtain significant direct CP asymmetry, one requires the
two interfering amplitudes to be of same order and their relative strong phase should be
significantly large (i.e., close to $\pi/2$). However, in the SM, the ratio of the CKM matrix
elements of the two terms in Eq. (20) can be given (in the Wolfenstein parametrization) as

$$|\lambda_u/\lambda_t| \simeq \lambda^2 \sqrt{\rho^2 + \eta^2} \simeq 2\%.$$  

Therefore, the first amplitude will be highly suppressed with
respect to the second unless $A_u >> A_t$. Hence, the naive expectation is that the direct CP
violation in the SM in this mode will be negligibly small. Using the generalized factorization
approach, we find $r = 0.15$, $\delta_t - \delta_u \sim 7^\circ$ and the direct CP violation in the mode $B^+ \rightarrow f_0K^+$
as of $\sim (-4\%)$. This in turn makes the mode interesting to look for the NP in terms of large
direct CP asymmetry.

In the presence of new physics the amplitude can be written as

$$A(B^+ \rightarrow f_0K^+) = A_{\text{SM}} + A_{\text{NP}} = A_{\text{SM}} \left[1 + r_{\text{NP}} e^{i\phi_{\text{NP}}} \right], \quad (22)$$
where $r_{NP} = |A_{NP}/A_{SM}|$, ($A_{SM}$ and $A_{NP}$ correspond to the SM and NP contributions to the $B^+ \rightarrow f_0 K^+$ decay amplitude, respectively) and $\phi_{NP} = \arg(A_{NP}/A_{SM})$, which contains both strong and weak phase components. The branching ratio for the $B^+ \rightarrow f_0 K^+$ decay process can be given as

$$BR(B^+ \rightarrow f_0 K^+) = BR^{SM} \left( 1 + r_{NP}^2 + 2r_{NP} \cos \phi_{NP} \right),$$

where $BR^{SM}$ represents the corresponding standard model value.

Now, we will present the basic formula of CP asymmetry parameters in the presence of new physics. Due to the contributions from new physics, these parameters deviate substantially from their standard model values. To find out the CP asymmetry, it is necessary to represent explicitly the strong and weak phases of the SM as well as of NP amplitudes. Although, it is expected that the SM amplitude $\lambda_u A_u$ is highly suppressed with respect to its $\lambda_t A_t$ counterpart, for completeness we will keep this term for the evaluation of $A_{CP}$.

We denote the NP contribution to the decay amplitude as $A_{NP} = |A_{NP}|e^{i\delta_n + \theta_n}$, where $\delta_n$ and $\theta_n$ denote the strong and weak phases of the NP amplitude, respectively. Thus, in the presence of NP, we can explicitly write the decay amplitude for $B^+ \rightarrow f_0 K^+$ mode as

$$A(B^+ \rightarrow f_0 K^+) = \lambda_u^*|A_u|e^{i\delta_u} + \lambda_t^*|A_t|e^{i\delta_t} + |A_{NP}|e^{i(\delta_n + \theta_n)}.$$  \hspace{1cm} (24)

The amplitude for $B^- \rightarrow f_0 K^-$ mode is obtained by changing the sign of the weak phases of the amplitude (24). Thus, the CP asymmetry parameter is given as

$$A_{CP} = \frac{2 \left( r \sin \gamma \sin \delta_{ut} + r_N \sin \theta_n \sin \delta_{nt} - rr_N \sin(\gamma - \theta_n) \sin \delta_{un} \right)}{|A|^2 - 2 \left( r \cos \gamma \cos \delta_{ut} + r_N \cos \theta_n \cos \delta_{nt} - rr_N \cos(\gamma - \theta_n) \cos \delta_{un} \right)}},$$

where $|A|^2 = 1 + r^2 + r_N^2$, $r_N = |A_{NP}/\lambda_t A_t|$ and $\delta_{ij} = \delta_i - \delta_j$ are the relative strong phases between different amplitudes.

Now we consider the CP violation parameters in the neutral $B$ meson decays, which has both direct and mixing-induced components. Let us consider the $B^0$ and $\bar{B}^0$ decay into a CP eigenstate $f_{CP}$ (we consider $f_{CP} = f_0 K_S$ with CP eigenvalue +1).

The time dependent CP asymmetry for $B \rightarrow f_0 K_S$ can be described as

$$A_{f_0 K_S}(t) = \frac{\Gamma(B^0(t) \rightarrow f_0 K_S) - \Gamma(\bar{B}^0(t) \rightarrow f_0 K_S)}{\Gamma(B^0(t) \rightarrow f_0 K_S) + \Gamma(\bar{B}^0(t) \rightarrow f_0 K_S)} = C_{f_0 K_S} \cos(\Delta M_{B_d} t) - S_{f_0 K_S} \sin(\Delta M_{B_d} t),$$

where $C_{f_0 K_S} = \frac{\Gamma(B^0(t) \rightarrow f_0 K_S)}{\Gamma(B^0(t) \rightarrow f_0 K_S) + \Gamma(\bar{B}^0(t) \rightarrow f_0 K_S)}$ and $S_{f_0 K_S} = \frac{-\Gamma(\bar{B}^0(t) \rightarrow f_0 K_S)}{\Gamma(B^0(t) \rightarrow f_0 K_S) + \Gamma(\bar{B}^0(t) \rightarrow f_0 K_S)}$.
where we identify
\[ C_{f_0 K_S} = \frac{1 - |\lambda|^2}{1 + |\lambda|^2} \quad \text{and} \quad S_{f_0 K_S} = \frac{2 \text{Im}(\lambda)}{1 + |\lambda|^2}, \]
(27)
as the direct and the mixing-induced CP asymmetries. The parameter \( \lambda \) corresponds to
\[ \lambda = \frac{q \mathcal{A}(\bar{B}^0 \to f_0 K_S)}{p \mathcal{A}(B^0 \to f_0 K_S)}, \]
(28)
where, \( q \) and \( p \) are the mixing parameters and are represented by the CKM elements in the standard model as
\[ \frac{q}{p} = \frac{V_{tb}^* V_{td}}{V_{tb} V_{td}^*} \sim \exp(-2i\beta). \]
(29)
Now the amplitude for \( \bar{B}^0 \to f_0 K_S \) can be symbolically written as
\[ \mathcal{A}(\bar{B}^0 \to f_0 K_S) = \lambda_t A_t, \]
(30)
where \( \lambda_t = V_{tb} V_{ts}^* \), which is real in the SM. Thus, the mixing induced CP asymmetry is given as,
\[ S_{f_0 K_S} = -\sin 2\beta, \]
same in magnitude as the one for \( B \to \psi K_S \), but with opposite sign and the direct CP asymmetry turns out to be identically zero.

However, the decay amplitude also receives some contribution from the internal up and charm quarks in the loop. Therefore, the CP violating parameters may deviate from their expected values. Now including the effects of \( u, c, t \) quarks in the loop and using CKM unitarity \((\lambda_u + \lambda_c + \lambda_t = 0)\), one can write the decay amplitude as
\[ \mathcal{A}(B^0 \to f_0 K^0) = \lambda_u^* A_u + \lambda_c^* A_c = \lambda_c^* A_c \left[ 1 + r'e^{i(\delta' + \gamma)} \right], \]
(31)
where the amplitude \( A_u \) contains contributions from \( u \) and \( t \) quarks in the loop (i.e., \( A_u = P_u - P_t \), where \( P_{u,c,t} \) are the penguin amplitudes corresponding to \( u, c, t \) quark exchange in the loop) and same argument holds for \( A_c \). The parameter \( r' \) is the ratio of the two amplitudes, i.e., \( r' = |\lambda_u A_u/\lambda_c A_c| \), \( \delta' = \delta_u - \delta_c = \text{Arg}(A_u/A_c) \) is the relative strong phase between them and \( \gamma \) is the weak phase. The explicit expressions for these amplitudes (in units of \(-G_F/\sqrt{2}\)) are given as
\[ A_q = \left[ a_0^q - \frac{a_{10}^q}{2} - r_{x_1} (a_0^q - \frac{a_8^q}{2}) \right] X + (2a_0^q - a_8^q) Y - \left[ a_0^q - \frac{a_{10}^q}{2} - (2a_0^q - a_8^q)m_B^2 \right] Z, \]
(32)
with \( q = u \) and \( c \), \( Y \) and \( Z \) are given in Eq. (12). Thus, one obtains the CP asymmetries as
\[ S_{f_0 K_S} = -\frac{\sin 2\beta + 2r' \cos \delta' \sin(2\beta + \gamma) + r'^2 \sin(2\beta + 2\gamma)}{1 + r'^2 + 2r' \cos \delta' \cos \gamma}, \]
\[ C_{f_0 K_S} = \frac{-2r' \sin \delta' \sin \gamma}{1 + r'^2 + 2r' \cos \delta' \cos \gamma}. \]
(33)
In order to know the precise value of the CP violating asymmetries one should know the values of \( r' \) and \( \delta' \). Using the QCD coefficients from [21] we obtain \( r' = 0.02 \), \( \delta' = 12^\circ \) and hence the CP asymmetries as

\[
S_{f_0 K_S} = -0.672 \quad \text{and} \quad C_{f_0 K_S} = -0.007 ,
\]

(34)

which are in accordance with the results of top quark dominance in the penguin loop. Therefore, here onwards we will consider the SM amplitude for the \( B^0 \to f_0 K_S \) process to be dominated by the top quark penguin.

New physics could in principle contribute to both mixing and decay amplitudes. The new physics contribution to mixing is universal while it is non-universal and process dependent in the decay amplitudes. As the NP contributions to mixing phenomena is universal, it will still set \( S_{\psi K_S} = -S_{f_0 K_S} \). Therefore, to explain the deviation between \( (S_{\psi K_S} \text{ and } (\sin 2\beta)_{f_0 K_S} = -S_{f_0 K_S}) \), here we explore the NP effects only in the decay amplitudes. Thus, including the NP contributions, we can write the decay amplitude for \( B \to f_0 K \) process as

\[
A(B^0 \to f_0 K^0) = A_{SM} + A_{NP} = \lambda_t^* A_t \left[ 1 - r_N \ e^{i(\delta_{nt} + \theta_n)} \right] ,
\]

(35)

where \( r_N = |A_{NP}/\lambda_t A_t| \), \( \delta_{nt} \) and \( \theta_n \) are the relative strong and weak phases between the new physics contributions to the decay amplitude and that of the SM part. The negative sign before \( r_N \) in Eq. (35) arises because the weak phase \( \pi \) of \( \lambda_t^* \) has been factored out. Thus, one can then obtain the expressions for the CP asymmetries as

\[
S_{NP}^{NP} = -\frac{\sin 2\beta - 2r_N \cos \delta_{nt} \sin(2\beta + \theta_n) + r_N^2 \sin(2\beta + 2\theta_n)}{1 + r_N^2 - 2r_N \cos \delta_{nt} \cos \theta_n} ,
\]

(36)

and

\[
C_{NP} = \frac{2r_N \sin \delta_{nt} \sin \theta_n}{1 + r_N^2 - 2r_N \cos \delta_{nt} \cos \theta_n} .
\]

(37)

Having obtained the CP asymmetry parameters, in the presence of new physics, we now proceed to evaluate the same in the R-parity violating supersymmetric model.

### IV. CONTRIBUTION FROM R-PARITY VIOLATING SUPERSYMMETRIC MODEL

We now analyze the decay modes in the minimal supersymmetric model with R-parity violation. In the supersymmetric models there may be interactions which violate the baryon
number $B$ and the lepton number $L$ generically. The simultaneous presence of both $L$ and $B$ number violating operators induce rapid proton decay, which may contradict strict experimental bound. In order to keep the proton lifetime within experimental limit, one needs to impose additional symmetry beyond the SM gauge symmetry to force the unwanted baryon and lepton number violating interactions to vanish. In most cases this has been done by imposing a discrete symmetry, called R-parity defined as, $R_p = (-1)^{(3B+L+2S)}$, which is +1 for all particles and −1 for all superparticles. This symmetry not only forbids rapid proton decay, but also prevents single creation and annihilation of superparticles. However, this symmetry is ad-hoc in nature. There is no theoretical arguments in support of this discrete symmetry. Hence, it is interesting to see the phenomenological consequences of the breaking of R-parity in such a way that either $B$ or $L$ number is violated, both not simultaneously violated, thus avoiding rapid proton decay. Extensive studies have been done to look for the direct as well as indirect evidence of R-parity violation from different processes and to put constraints on various R-parity violating couplings.

For our purpose, we will consider the Lepton number violating super-potential with only $\lambda'$ couplings, which is given as

$$W' = \lambda'_{ijk} L_i Q_j D^c_k,$$  \hspace{1cm} (38)

where $i, j, k$ are generation indices, $L_i$ and $Q_j$ are $SU(2)$ doublet for lepton and quark superfields and $D^c_k$ is the down type quark singlet superfield.

Thus the effective Hamiltonian for charmless hadronic $B$ decays can be given as

$$H^X_{\text{eff}} = d^R_{jkn}[\bar{d}_n \gamma_\mu^L d_j \gamma_\mu^R b_k d_R j] + d^L_{jkn}[\bar{d}_n \gamma_\mu^L d_j \gamma_\mu^R b_k d_L j] + u^R_{jkn}[\bar{u}_k \gamma_\mu^L u_j \gamma_\mu^R b_n d_R j],$$  \hspace{1cm} (39)

where $\alpha, \beta$ are the color indices, $\gamma_\mu^{R,L} = \gamma^\mu(1 \pm \gamma_5)$ and

$$d^R_{jkn} = \sum_{i=1}^3 \frac{\lambda'_{ijk} \lambda^*_{jkn}}{8m^2_{\nu_{Li}}}, \quad d^L_{jkn} = \sum_{i=1}^3 \frac{\lambda'_{ijk} \lambda^*_{jkn}}{8m^2_{\nu_{Li}}}, \quad u^R_{jkn} = \sum_{i=1}^3 \frac{\lambda'_{ijm} \lambda^*_{knl}}{8m^2_{e_{Li}}}. \hspace{1cm} (40)$$

Thus one can write the transition amplitudes as

$$A^{X}(B^+ \rightarrow f_0 K^+) = -2(d^{L}_{222} + d^{R}_{222}) Y + u^{R}_{112} r_\chi X - (d^{R}_{112} + d^{L}_{121}) 2Y_d,$$

$$A^{X}(B^0 \rightarrow f_0 K^0) = -2(d^{L}_{222} + d^{R}_{222}) Y + (d^{R}_{121} + d^{L}_{112}) r_\chi X + (d^{R}_{112} + d^{L}_{121}) \left(\frac{X}{N} - 2Y_d\right),$$  \hspace{1cm} (41)
where $Y_d$ is the value of $Y$ with $\bar{f}_{f_0}^s$ replaced by $\bar{f}_{f_0}^d$.

Following the standard practice, we shall assume that the RPV couplings are hierarchical, i.e., only one combination of the coupling is numerically significant. Furthermore, we also assume that both the transitions $B^{+,0} \to f_0 K^{+,0}$ receive dominant contribution from the quark level transition $b \to s\bar{s}s$, and hence we consider only $d_{222}^L$ coupling to be nonzero. As discussed in Ref. [31], we will also discard the $d_{222}^R$ coupling in our analysis, as it is related to $u_{222}^R$ by SU(2) isospin symmetry and its effect in the mode $B \to J/\psi K_S$ is found to be negligibly small. Thus, with these approximations the transition amplitudes for both the processes can be given as

$$A^\lambda(B \to f_0 K) = -\frac{1}{8m_{\tilde{\nu}^i_{Li}}} \left(\lambda'_{i32} \lambda'^*_{i22}\right) 2m_{f_0} \bar{f}_{f_0}^s \frac{m_{f_0} - m_{K}}{m_{b} - m_{s}} F_{BK}(q^2),$$

where the summation over $i = 1, 2, 3$ is implied. Now considering the values of R-parity couplings as

$$\lambda'_{i32} \lambda'^*_i = Re^{i\theta_n},$$

where $R = |\lambda'_{i32} \lambda'^*_i|$ and $\theta_n$ is the new weak phase with range $-\pi \leq \theta_n \leq \pi$. It should be noted that since the dominant SM amplitude (i.e., the t-quark dominated penguin amplitude $A_t$) contains the weak phase $\pi$, we vary the weak phase $\theta_n$ between $[-\pi, \pi]$, so that the NP amplitude will interfere constructively with the SM amplitude when the relative weak phase between them is zero. To see the effect of R-parity violation in the decay modes $B \to f_0(980)K$, it is essential to know the value of the RPV couplings ($R$). We first present a crude estimation of $R$ by assuming that R-parity will explain the observed discrepancy between the observed and SM predicted branching ratios for $B \to f_0 K$ modes. We will further assume that the new physics amplitude will interfere constructively with the standard model amplitude (i.e., $\phi_{NP} = 0$ in Eq. (22)), so that one can obtain a lower bound on $r_{NP}$ from Eqn. (23). Now using the values of the experimental branching ratios from Eq. (4) and the corresponding SM values from (19), we obtain the lower bound as $r_{NP} \geq 0.6$. This, in turn with Eqns. (11), (14) and (42) gives

$$R \geq 1 \times 10^{-3},$$

for $m_{\tilde{\nu}^i_{Li}} = 100$ GeV. Recently, in Ref. [32] it has been shown that the branching ratio and the polarization anomaly in $B \to \phi K^*$ modes can be resolved in the R-parity violating supersymmetric model for a very narrow interval in the parameter space as $|\lambda'_{i32} \lambda'^*_i|/m_{\tilde{\nu}^i_{Li}}^2 \in$
[1.5 \times 10^{-3}, 2.1 \times 10^{-3}]$, for the sneutrino mass scale $m_{\tilde{\nu}_{Li}} = 100$ GeV. Therefore, in this analysis we consider the lowest value for $R$ i.e., $1.5 \times 10^{-3}$ from the above allowed range, which also satisfies the constraint (44). Using this value we obtain the ratios of RPV to SM amplitude, as defined in section III, as

\begin{align*}
    r_{NP} &= 0.81, \quad r_N = 0.87 \quad \text{(for $B^+ \to f_0 K^+$)} ,
    \\
    r_N &= 0.92, \quad \text{(for $B^0 \to f_0 K^0$)}.
\end{align*}

Therefore, the upper limits in the branching ratios (for $\phi_{NP} = 0$ in Eq.(23)) in the RPV model are found to be

\begin{align*}
    BR(B^+ \to f_0(980) K^+) &\leq 21.6 \times 10^{-6} ,
    \\
    BR(B^0 \to f_0(980) K^0) &\leq 17.4 \times 10^{-6} .
\end{align*}

Thus one can see that the observed branching ratios (4) can be accommodated in the RPV model.

Now assuming the strong phase difference $\delta_{nt}$ to be small (e.g., $\sim 10^\circ$), direct CP violation for $B^+ \to f_0 K^+$ process and the mixing induced CP violating parameter from $B^0 \to f_0 K^0$ are shown in figures - 3 and 4. Thus, as seen from the figures, the observed $(\sin 2\beta)_{f_0 K_S} = -S_{f_0 K_S} = 0.51 \pm 0.19$ can be explained in the RPV model and large direct CP violation (upto 80 %) in $B^+ \to f_0 K^+$ mode could be obtainable in this model. However, there is no obvious reason why the strong phase difference $\delta_{nt}$ could be small. To see the impact of the strong phase we vary it between the range of $-\pi$ and $\pi$ and plot the correlation between direct and mixing induced CP asymmetries for $B^0 \to f_0 K_S$, for two representative values of weak phase $\theta_n = \pi/2$ and $\pi/4$, in figure-5. From the figure, it is seen that R-parity violating supersymmetric model can accommodate large CP violation in the $B^0 \to f_0 K_S$ decay mode.

V. CONCLUSION

In this paper we have studied the rare decay modes $B \to f_0(980) K$, involving a scalar and a pseudoscalar meson in the final state. Since the structure of the $f_0$ meson is not well established till now, we consider it as a $q\bar{q}$ state, comprising of both $s\bar{s}$ and $(u\bar{u} + d\bar{d})/\sqrt{2}$ components with a mixing angle of 138°, which appears to be the most preferable one. Using the generalized factorization approach, we found that the branching ratios in the standard
FIG. 3: Direct CP violation for the process $B^+ \rightarrow f_0(980)K^+$ versus the new weak phase $\theta_n$ in degrees.

FIG. 4: Mixing induced CP violation for the process $B^0 \rightarrow f_0(980)K^0$ versus the new weak phase $\theta_n$ in degrees.

model are below the current experimental values, as was obtained in previous studies using different approaches. The average value of the observed mixing induced CP asymmetry, i.e., $(\sin 2\beta)_{f_0K_S} = -S_{f_0K_S}$, also has about one sigma deviation from that of $S_{J/\psi K}$. To explain the observed discrepancy in the branching ratios and CP asymmetry parameter, we considered the R-parity violating model. Since these processes receive dominant contribution from $b \rightarrow s\bar{s}s$ loop induced penguins, we assumed that the new physics parameters will affect such transitions strongly. We found that the R-parity violating model can explain the observed discrepancy in the branching ratios and the CP violation parameter $-S_{f_0K_S}$. It can accommodate large CP violation even for small relative strong phase between SM and RPV amplitudes. In this analysis, we have considered a representative value for the RPV coupling $|\lambda_{t32}'\lambda_{t22}^*|$. But, it should be noted that using the data on the branching ratios and
FIG. 5: The correlation plot between $S^{NP}$ and $C^{NP}$ for the process $B^0 \rightarrow f_0(980)K_S$ in the RPV model for two representative values of weak phases ($\theta_n = \pi/2, \pi/4$), where we have used $r_N = 0.92$, and varied the strong phase $\delta_{nt}$ between $-\pi$ and $\pi$.

CP asymmetries of the processes, which have dominant $b \rightarrow s\bar{s}s$ quark level transitions, it would be possible to obtain the allowed parameter space for the magnitudes and phases of the RPV couplings. If, in future, the $q\bar{q}$ structure for $f_0$ is established then these modes could also play an important role to look for new physics beyond the standard model or else, at least, it will certainly enrich our understanding regarding the nature of the light scalar mesons.

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