Some Rare Decays of Particles within R-parity violating Super-symmetric Model

By

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CIIT/FA04-PPH-002/ISB

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Declaration

I Azeem Mir hereby declare that I have produced the work presented in this thesis, during the scheduled period of study. I also declare that I have not taken any material from any source except referred to wherever due. If a violation of HEC rules on research has occurred in this thesis, I shall be liable to punishable action under the plagiarism rules of the HEC.

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It is certified that Mr. Azeem Mir has carried out all the work related to this thesis under my supervision at the Department of Physics, CIIT Islamabad.

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Dedicated

To

My Sister and Parents

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Abstract

Some Rare Decays of Particles within R-parity violating Super-symmetric Model

This thesis compares experimental studies of two and three body leptonic decays of mesons (***B*** and ***K***) with theoretical predictions of R-parity violating Minimal Supersymmetric Standard Model. Observables like branching fraction, forward backward asymmetry and polarization asymmetry are studied regarding these decays. Forward backward asymmetry is found to be significant in the case of $K^{\pm} \to \pi^{\pm}\mu^{+}\mu^{-}$ only and vanishingly small in case of $B^{\pm} \to K^{\pm}\mu^{+}\mu^{-}$. Polarization asymmetry has also been studied in this research work. It is found to be significant only for two body leptonic decays and comparable with standard model in case of three body leptonic decay of mesons. An error analysis of the branching fraction and CP-asymmetry has also been made for three body leptonic decay of mesons. Theoretical predictions for both branching fraction and CP-asymmetry agree well with present experimental data.

List of thesis Papers

This thesis is based on the following paper, whose copies are attached at the end.

1. Implications of R parity violating Yukawa couplings in DeltaS=1 semileptonic decays of K mesons.
   Phys.Rev.D76:114005, 2007.
   (A. Mir, F. Tahir, Mahnaz Q. Haseeb, K. Ahmed)

2. Lepton polarization asymmetry in B -->l+l- decays in R-parity - violating minimal super-symmetric standard model.
   Europhys.Lett.83:41002, 2008.
   (A. Mir, F. Tahir, K. Ahmed)

3. Probing new physics through B ---> K l+ l- decays in R-parity violating minimal Super-symmetric standard model.
   Eur.Phys.J.C65:475-490, 2010.
   (A. Mir, F. Tahir, K. Ahmed)

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Chapter 1

Introduction

Standard model (SM) was introduced by Sheldon Glashow, Steven Weinberg and Abdus Salam [1], based on the Yang and Mills unified gauge theory [2], which was extended to incorporate higher order corrections by G.'t Hooft and M. Veltman [3]. It gives a coherent framework describing the world of particle physics. SM relates fundamental building blocks (particles and fields) in a qualitative and quantitative manner. It has also been proven to be accurate and precise to all laboratory tests up-to-date. This can be demonstrated by comparing the experimental measurement of $W^{\pm}$, Z^0 and top quark mass with the predictions of SM given in table. (1.5.1) [4].

SM gives us a framework to study all known particles and fields except gravity. The framework consists of gauge interactions (electroweak and strong) present among all known elementary particles. One missing ingredient of this framework is the Higgs particle, which is theorized to give mass to bosons and fermions while preserving gauge invariance [5]. Higgs boson introduces many other problems. Virtual particles of the SM in the vacuum contribute to Higgs mass in such a way that mass of Higgs exceeds 1 TeV scale [6]. One must then discover a new mechanism of electroweak symmetry breaking or find a new symmetry to prevent the large contribution to Higgs mass [6].

Supersymmetry (SUSY) is one such symmetry that postulates new particles (matter constituents and bosons) called sparticles [7]. Each sparticle cancels the large contribution contributed to Higgs mass by its partner particle in exact SUSY limit. Since the spartners do not have the same mass as that of their SM partners, therefore, SUSY must be broken [8].

SUSY also helps in bringing strong electroweak unification. In SU(5) SUSY formalism, coupling strengths of strong and electroweak interactions become the same at $10^{16} GeV$ [9]. This behavior is shown to be extrapolated upto $10^{19} GeV$. SUSY plays a major role in particle physics phenomenology. It may give insight into the origin of neutrino mass. Particle processes that violate baryon and lepton number are possible within the framework of SUSY.

The superpotential of Minimal supersymmetry that assumes minimum number of particles and fields, is divided into two parts [7], namely R-parity conserving and R-parity vioalting respectively. This is mainly due to phenomenological reasons i.e. the absence of evidence of lepton and baryon number violating processes. R-parity conserving excludes such processes while R-parity violating admits such processes. Most importantly, flavor changing neutral currents (FCNC) are possible at both loop and tree level in R-parity violating SUSY [10].

FCNC involve those processes, in which flavor of fermions is changed without the change of charge. Examples of such processes include leptonic decay of hadrons, oscillation of neutral mesons and top quark decays. Leptonic decay of hadrons are suppressed in SM as they are forbidden at tree level. Only higher order Feynman diagrams like box and penguin contribute to these processes. FCNC may reveal the effect of virtual particles (present in loop diagrams) associated with new physics without directly producing them [10]. Therefore, FCNC provide a window to look for physics beyond SM.

The chapter-wise plan is described as follows: In section 1.1, a concise framework of SM is reviewed. Section 1.1(a) deals with the electroweak unification. Higgs mechanism is reviewed in section 1.1(b), while Two Higgs Doublet Model is reviewed in section 1.1(c). In section 1.2 problems with SM are briefly introduced with main emphasis on Higgs hierarchy. Supersymmetry, Minimal Supersymmetric SM (MSSM), Superpotential, Hierarchy problem and R-parity conservation are reviewed in sections 1.3. In section 1.4, R-parity violation and its phenomenology is discussed. Section 1.5 discusses collider (*LEP*, *HERA*, *Tevatron* and *LHC*) searches for R-parity violation.

In chapter 2, we review and discuss the role of R-parity violation in semileptonic charged kaon decays. Section 2.1 introduces semileptonic charged kaon decays. In the context, the contribution to flavor changing neutral currents coming from R-parity violating interactions are elaborated here. In section 2.2, the framework for the analysis is presented. Section 2.3

discusses the results of our analysis.

In chapter 3, we discuss the role of R-parity violation in pure leptonic decays of neutral mesons (K and B). Section 3.1 introduces leptonic decays of neutral K and B mesons briefly. In section 3.2, the framework for the analysis is given. In Section 3.3, the results of our analysis are discussed.

In chapter 4, we discuss the role of R-parity violation in semileptonic decays of charged B mesons. Section 4.1 introduces semileptonic decays of B mesons. It also gives a brief overview of previous work done by several authors. In section 4.2, the framework of our analysis is presented. In Section 4.3, we discuss the results obtained from our calculations.

In chapter 5, we present the conclusion and discuss future work.

1.1 Standard Model

SM consists of three fundamental forces of nature and 12 elementary particles. The electromagnetic force (carried by photon) is exchanged between two charged particles. The coupling strength α of electromagnetic field can be given by [6]

$$\alpha = \frac{e^2}{4\pi\varepsilon_o \hbar c} \approx \frac{1}{137}.$$

The free Lagrangian for electromagnetic interaction is given by [6]

$$\mathcal{L} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} - J^\rho A_\rho, \tag{1.1.1}$$

where $F_{\mu\nu}$ is given by

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu.$$

A_μ is the four-vector potential and J_μ is the four-vector current given by

$$A^\mu = (V, \vec{A});\ \ J^\mu = (\rho, \vec{J}).$$

The complete gauge invariant Lagrangian including fermionic terms is given as

$$\mathcal{L} = \bar{\psi}(i\gamma^{\mu}D_{\mu} - m)\psi - \frac{1}{4}F^{\mu\nu}F_{\mu\nu} - J^{\rho}A_{\rho}, \tag{1.1.2}$$

where

$$D_{\mu} = \partial_{\mu} - ieA_{\mu}.$$

The weak nuclear force was discovered during the investigation of nuclear beta decay, e.g.

$$n \rightarrow p + e^{-} + \bar{\nu}_e.$$

Its strength is given by G_F, where

$$G_F = 1.166 \times 10^{-5} GeV^{-2}.$$

$W^{\pm}$ bosons affect left handed leptons (doublets) only. Z^0 bosons affect both left and right handed leptons (except for the neutrinos, which have no right handed component under SM gauge group) [6].

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L, \tag{1.1.3}$$

$$e_R, \mu_R, \tau_R. \tag{1.1.4}$$

The free Lagrangian for weak interactions is given by

$$\mathcal{L} = -\frac{1}{4}W_a^{\mu\nu}W_{\mu\nu}^a \tag{1.1.5}$$

Where

$$W_{\mu v}^a = \partial_\mu W_v^a - \partial_\nu W_\mu^a + g \sum_{b,c} \varepsilon_{abc} W_\mu^b \times W_v^c, \tag{1.1.6}$$

W_a^{μ} is the gauge field of $SU(2)_L$. It couples to the fermion via coupling strength g. It can be identified with the gauge bosons $(W_\mu^1, W_\mu^2, W_\mu^3)$. The third term in eq. (1.1.6) is due to the non-abelian character of the group. The physical bosons ($W^{\pm}$) are given as [6]

$$\begin{aligned} W_\mu^+ &= (W_\mu^1 - iW_\mu^2)/\sqrt{2}, \\ W_\mu^- &= (W_\mu^1 + iW_\mu^2)/\sqrt{2}, \\ W_\mu^0 &= W_\mu^3. \end{aligned}$$

Here, W_μ^0 does not correspond to physical Z^0 boson. This is because weak neutral current has both left and right handed components unlike W_μ^0, which is the gauge field of $SU(2)_L$. The complete Lagrangian including fermionic terms is given by

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu D_\mu - m)\psi - \frac{1}{4}W_a^{\mu\nu}W_{\mu\nu}^a, \tag{1.1.7}$$

where the covariant derivative:

$$D_\mu = \partial_\mu - ig\frac{\tau^k}{2}W_\mu^k - ig'B_\mu \tag{1.1.8}$$

The free Lagrangian for strong interaction is given by [6]

$$\mathcal{L} = -\frac{1}{4}G_a^{\mu\nu}G_{\mu\nu}^a,$$

where

$$G_a^{\mu v} = \partial^\mu G_a^v - \partial^v G_a^\mu - g_s f_{abc} G_\mu^b G_\nu^c.$$

The f_{abc} are structure constants; for $SU(3)$. The complete lagrangian including fermionic terms is given by [6]

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu D_\mu - m)\psi - \frac{1}{4}G_a^{\mu\nu}G_{\mu\nu}^a, \tag{1.1.9}$$

where the covariant derivative

$$D_\mu = \partial_\mu - ig_s\frac{\lambda^b}{2}G_\mu^b.$$

Hadrons (mesons and baryons) are made up of constituent particles called quarks. Quarks carry color charge (red, green and blue), whereas antiquarks carry anti-color charges. Gluons carry a color and an anti-color. These gluons are responsible for interaction between quarks and thus interaction between hadrons. In the framework of Quantum Chromodynamics (QCD), there are nine possible color-anticolor combinations representing a gluon. One of these states is a color singlet, given as [6]

$$(r\bar{r} + b\bar{b} + g\bar{g})/\sqrt{3}.$$

This color singlet state does not exist physically. Remaining eight of these possible nine color-anticolor combinations exist physically. These eight states form a "color octet".

$$(r\bar{b} + b\bar{r})/\sqrt{2},$$
$$i(b\bar{r} - r\bar{b})/\sqrt{2},$$
$$(r\bar{g} + g\bar{r})/\sqrt{2},$$
$$i(g\bar{r} - r\bar{g})/\sqrt{2},$$
$$(b\bar{g} + g\bar{b})/\sqrt{2},$$
$$i(g\bar{b} - b\bar{g})/\sqrt{2},$$
$$(r\bar{r} - b\bar{b})/\sqrt{2},$$
$$(r\bar{r} + b\bar{b} - 2g\bar{g})/\sqrt{6}.$$

(a) Electroweak Unification

Electromagnetic and weak interactions are different from each other at ordinary energies. As electromagnetic interaction is described by an abelian gauge group, while the weak interactions are described by the non-abelian gauge group. Electromagnetic interaction is long range, while weak interactions are of short range ($10^{-18}m$).

The weak neutral current has a right handed component along with left handed component. Since the weak interaction lagrangian (see eq. (1.1.5)) is invariant under $SU(2)_L$ transformations only, so the neutral current cannot be completely described by such lagrangian. Some kind of right handed gauge group transformations must exist along with $SU(2)_L$. The gauge group for Electromagnetic interaction $U(1)_{e.m}$ is one such choice [6]. Thus a lagrangian invariant under the $SU(2)_L \times U(1)_Y$ transformation may describe neutral currents. Such a unification scheme was first proposed by Glashow in 1961, long before the discovery of the weak neutral current and then was extended by Weinberg and Salam to include the massive vector bosons

($W^{\pm}$ and Z^0) [1]. Here, Y is the weak hypercharge given by [5].

$$Y = Q - \frac{T_3}{2}.$$

The hypercharge Y is the generator of $U(1)_Y$. There are two types of hypercharges Y_L and Y_R, defined in the gauge transformation of the doublet ψ_L (given by eq. (1.1.3)) and singlet ψ_R (given by eq. (1.1.4)) respectively.

$$\begin{aligned}\psi'_L &= e^{-iY_L\alpha(x)}\psi_L, \quad SU(2)_L \\ \psi'_R &= e^{-iY_R\alpha(x)}\psi_R. \quad U(1)_R\end{aligned}$$

The lagrangian for $SU(2)_L \otimes U(1)_Y$ is given by [1, 6]:

$$\mathcal{L} = \bar{L}\gamma^\mu D_\mu L + \bar{R}\gamma^\mu D'_\mu R - \frac{1}{4}B^{\mu\nu}B_{\mu\nu} - \frac{1}{4}W_a^{\mu\nu}W^a_{\mu\nu} + (D_\mu\phi)^\dagger(D^\mu\phi) - V(\phi,\phi^\dagger) - g_1\bar{L}\phi R + g_2\bar{L}\widetilde{\phi}R \tag{1.1.10}$$

where the terms containing $\phi, \phi^\dagger$ describe Higgs scalar boson and their interaction with fermions,

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu.$$

B_μ is the gauge field of $U(1)_Y$. It couples with the fermion via coupling strength g'. The covariant derivatives D_μ and D'_μ are now given as

$$D_\mu = \partial_\mu - ig'B_\mu - ig\frac{\tau^k}{2}W^k_\mu; \;\; D'_\mu = \partial_\mu - ig'B_\mu;$$

The physical fields A_μ (electromagnetic current) and Z_μ (neutral current) are given in terms of gauge fields B_μ and Z_μ [6]:

$$\begin{aligned}A_\mu &= B_\mu\cos\theta_W + W^3_\mu\sin\theta_W, \\ Z_\mu &= -B_\mu\sin\theta_W + W^3_\mu\cos\theta_W.\end{aligned}$$

θ_W is the Weinberg weak mixing angle. Electrical charge (e) is related to weak couplings strength (g) and gauge coupling strength $g'(B_\mu)$ by the following relation

$$e = g\sin\theta_W = g'\cos\theta_W.$$

(b) Higgs and Spontaneous Symmetry Breaking

The idea of Higgs boson arose to answer the lack of gauge invariance in the free lagrangian of electro-weak interaction (see eq. (1.1.10)) [5]. Weak interactions are actually mediated by massive W and Z bosons. The mass term in the interaction Lagrangian eq. (1.1.7) makes it non-renormalizable [6]. In contrast, photons are massless, so the free lagrangian for photon field is automatically gauge invariant.

Also there must be a symmetry ($SU(2)_L$) between electron and electron-neutrino (ν_e), which is not possible due to the difference between the masses of neutrino and electron. Further, it can be demonstrated that it is possible to construct a left handed gauge theory from the $SU(2)_L$ transformations only in the case for mass-less fermions [6]. These considerations make it evident that there must be some mechanism, which respects not only the gauge invariance and $SU(2)_L$ symmetry but also provides masses to the fermions and bosons. Spontaneous symmetry breaking is one such simplest mechanism present in nature like in condensed matter physics (Curie Temperature etc) [6].

Spontaneous symmetry breaking occurs whenever a system that obeys some symmetry group makes a transition into a non-symmetric vacuum state. Examples include from ferromagnetic material to breaking of translational symmetry in crystal lattice [6]. In SM, this is achieved by introducing Higgs boson. The Higgs boson provides mass to SM particles via interaction. It does so without disturbing the gauge invariance of Lagrangian of $W^\pm$ and Z bosons. Further, Higgs boson must be neutral and scalar (spin -0). The Higgs potential is given by [5]

$$\begin{aligned} V(\phi^\dagger, \phi) &= \mu^2\phi^\dagger\phi + \lambda(\phi^\dagger\phi)^2 \\ \mu^2 &< 0;\ \lambda > 0, \end{aligned}$$

where μ represents the mass of scalar particle. The lagrangian density $\mathcal{L}_\phi$ for ϕ is then given as

$$\mathcal{L}_\phi = (\partial_\rho\phi)^\dagger(\partial^\rho\phi) - (\mu^2\phi^\dagger\phi + \lambda(\phi^\dagger\phi)^2). \tag{1.1.11}$$

It can be shown that the potential energy is minimum for

$$\phi = \pm \text{v}, \text{ where v}=\sqrt{-\mu^2/\lambda}.$$

The $SU(2)_L$ invariant lagrangian density $\mathcal{L}_\phi$ for ϕ is then given as

$$\mathcal{L}_\phi = (D_\rho \phi)^\dagger (D^\rho \phi) - (\mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2),$$

where the covariant derivative D_ρ is given by eq. (1.1.8). The Higgs field ϕ is now associated with an $SU(2)_L$ doublet. A choice may be [5]

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}, \tag{1.1.12}$$

where ϕ^+ and ϕ^0 are each complex fields,

$$\begin{aligned} \phi^+ &= \frac{\phi_1 + i\phi_2}{\sqrt{2}}, \\ \phi^0 &= \frac{\phi_3 + i\phi_4}{\sqrt{2}}, \\ \phi^\dagger \phi &= \frac{\phi_1^2 + \phi_2^2 + \phi_3^2 + \phi_4^2}{2} \\ \phi^\dagger \phi &= -\frac{\mu^2}{2\lambda} = \frac{\text{v}^2}{2}. \end{aligned} \tag{1.1.13}$$

Eq. (1.1.13) therefore, suggests that there are many ways to have eq. (1.1.12). One appropriate and simple choice is [5]

$$\phi_{\text{min}} = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \text{v} \end{pmatrix}, \text{ where v}=\sqrt{-\mu^2/\lambda}.$$

Fermion masses are generated through the coupling of the left and right handed fermions to a scalar doublet Higgs boson ϕ [5, 6]

$$\mathcal{L}_{Yukawa} = -h_f [\bar{\psi}_L \phi \psi_R + \bar{\psi}_R \phi^\dagger \psi_L] + h.c. \tag{1.1.14}$$

Fermion masses m_f are given as:

$$m_f = h_f \text{v},$$

where, ψ_L and ψ_R are defined in section 1.1(a) and h_f is the coupling of fermions with the Higgs boson. The W boson mass is given by:

$$M_W = \frac{1}{\sqrt{2}} g \text{v},$$

where

$$\text{v} = (\frac{1}{\sqrt{2} G_F})^{1/2} = 246 \ GeV. \qquad (1.1.15)$$

From Higgs mechanism, one can derive

$$M_Z = \frac{M_W}{\cos\theta}.$$

The physical Higgs boson mass is given by [5, 6]:

$$M_H = \text{v}\sqrt{2\lambda}.$$

Both the lower and upper bound on Higgs mass can be derived from the considerations of vacuum stability and triviality [11].

$$(114 < M_H < 185) \ GeV/c^2.$$

Recent data further excludes the Higgs mass between (**159** and **168**) GeV/c^2 with a unitarity limit of 1TeV [5].

(c) Two Higgs Doublet Model

In nature there may be more than one Higgs bosons [12]. This is due to the fact that there are many other ways (any model with arbitrary number of singlet and doublet Higgs bosons) to satisfy [6].

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} \approx 1.$$

$FCNC$ is not allowed at the tree level by additional Higgs doublets if the given model satisfies following conditions:

(1) Arranging the Higgs masses predicted by the model to be so heavy that the tree level $FCNC$ are suppressed.

(2) The fermions of a given charge couple to no more than one Higgs doublet. It was proved by Glashow and Weinberg [13] that minimal SUSY contains two Higgs doublets satisfying this particular condition.

Two Higgs doublet model (2HDM) is one such minimal extension of the Higgs sector. It is minimal in the sense that few new parameters are added. Two complex scalar fields $\phi_1$and ϕ_2, associated with $SU(2)_L$ doublet are introduced in this model. The Higgs potential is then given by [5]

$$\begin{aligned} V(\phi_1,\phi_2) &= \lambda_1(\phi_1^\dagger\phi_1 - \mathrm{v}_1^2) + \lambda_2(\phi_2^\dagger\phi_2 - \mathrm{v}_2^2) + \lambda_3[(\phi_1^\dagger\phi_1 - \mathrm{v}_1^2) + (\phi_2^\dagger\phi_2 - \mathrm{v}_2^2)]^2 \\ &+\lambda_4[(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) - (\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1)] + \lambda_5[\mathrm{Re}[\phi_1^\dagger\phi_2] - \mathrm{v}_1\mathrm{v}_2\cos\alpha]^2 \\ &+\lambda_5[\mathrm{Im}[\phi_1^\dagger\phi_2] - \mathrm{v}_1\mathrm{v}_2\sin\alpha]^2, \end{aligned}$$

where the $\lambda_i' s$ are all assumed to be real and positive. The minimum of this potential is given by [5]

$$<\phi_1>= \begin{pmatrix} 0 \\ \mathrm{v}_1 \end{pmatrix}; <\phi_2>= \begin{pmatrix} 0 \\ \mathrm{v}_2 e^{i\alpha} \end{pmatrix}.$$

Here, α is the phase factor. A parameter that relates the two expectation values v_1 and v_2 is β given as:

$$\tan\beta = \frac{\mathrm{v}_2}{\mathrm{v}_1}.$$

This model introduces five physical Higgs bosons, ($H^{\pm}$, H^0 and h^0 (CP-even) and A^0 (CPodd)). These physical states are given by [5];

$$\begin{aligned}
H^{\pm} &= -\phi_1^{\pm}\sin\beta + \phi_2^{\pm}\cos\beta \\
H^0 &= \sqrt{2}[(\mathrm{Re}\,\phi_1^0 - \mathrm{v}_1)\cos\gamma + (\mathrm{Re}\,\phi_2^0 - \mathrm{v}_2)\sin\gamma] \\
h^0 &= \sqrt{2}[-(\mathrm{Re}\,\phi_1^0 - \mathrm{v}_1)\sin\gamma + (\mathrm{Re}\,\phi_2^0 - \mathrm{v}_2)\cos\gamma] \\
A^0 &= \sqrt{2}(-\,\mathrm{Im}\,\phi_1^0\sin\beta + \mathrm{Im}\,\phi_2^0\cos\beta).
\end{aligned}$$

Here, γ mixes H^0 and h^0. The Goldstone boson states are given as:

$$\begin{aligned}
G^{\pm} &= \phi_1^{\pm}\cos\beta + \phi_2^{\pm}\sin\beta \\
G^0 &= \sqrt{2}(\mathrm{Im}\,\phi_1^0\cos\beta + \mathrm{Im}\,\phi_2^0\sin\beta).
\end{aligned}$$

These Goldstone bosons are then eaten up by the W and Z bosons. The W mass is then given by [5]

$$M_W = g\sqrt{\frac{\mathrm{v}_1^2 + \mathrm{v}_2^2}{2}}.$$

Summing up, the model predicts five higgs bosons with six free parameters. These six parameters include masses of four Higgs bosons, tanβ and Higgs mixing angle γ.

1.2 Physics beyond SM

SM has yet satisfied all current experimental data on electroweak and QCD interactions. However, SM has appears adhoc and incomplete as a theory. It cannot incorporate gravity. Attempts to explain the origin of neutrino mass are also not successful yet. SM further cannot explain cosmological phenomenas like the nature of dark matter and dark energy. Issues such as those provide need and motivation to go beyond SM. Then there is the Hierarchy problem related to the Higgs mass [6]. Pattern in Fermion masses have also been left unexplained by the SM [14]. Also, is the mystery of why leptons and quarks are replicated in exactly three families.

(a) Hierarchy problem

Hierarchy problem is a representation of the fact that the Higgs mass is extremely sensitive to new physics [15]. SM is a low energy effective theory i.e., it is a part of a large theory that is important at only higher energy scale. Quantum gravity is one such scale, which is effective at Planck's mass scale [6].

$$M_P = 1.22 \times 10^{19} GeV.$$

At such high energy scale, virtual particles of the vacuum give a divergent contribution to Higgs boson through Higgs self interaction (see Fig. (1.5.1 and 1.5.2)). Fig. (1.5.1) shows contribution of virtual scalar and fermions present in quantum vacuum to the Higgs mass. The contribution to Higgs mass at given energy scale Λ by fermions is given by [15]

$$(\mu^2_{phys})_f - \mu^2 \sim -\lambda_f^2 \ \Lambda^2 + \log \Lambda \text{terms}.$$

Similarly, the contribution to Higgs mass at given energy scale Λ by scalar particles is given by

$$(\mu^2_{phys})_S - \mu^2 \sim \lambda_S \ \Lambda^2 + \log \Lambda \text{ terms}.$$

At Max Planck energy scale i.e.,

$$\Lambda = M_P \sim 10^{19} GeV,$$

the one loop correction will be $\sim (10^{38} GeV^2)$. It implies that we have to use a very large value of μ to give a phenomenological acceptable μ_{phys}. This is called fine-tuning problem [15]. *SUSY* proposes an elegant solution. If we assume that each fermion has a scalar partner and their coupling strength to Higgs boson is same i.e.,

$$\lambda_S = \lambda_f^2 = \lambda,$$

the term Λ^2 will be absent in the total contribution to Higgs mass by the fermions and scalar particles in the vacuum i.e.,

$$(\mu^2_{phys})_f + (\mu^2_{phys})_S \sim \ \mu^2.$$

Since spartners have not been discovered yet, so spartners must be a lot heavier than SM particles. The quadratic divergences cancel out again but another term will survive i.e., [15]

$$\lambda(m^2_S - m^2_f) \ln \Lambda.$$

The factor $(m^2_S - m^2_f)$ must remain of the order of $\mathbf{v}^2$(see eq.(1.1.15)) i.e., $(246\, GeV)^2$, otherwise, fine tuning will be necessary again. Supersymmetry also proposes that bosons of SM must have fermionic spartners, so that their contributions to the Higgs boson is canceled out neatly (see Fig. (1.5.1 and 1.5.2)).

LHC is currently exploring at the scale of energies beyond $1TeV$ and it is most likely that if $SUSY$ is a viable model for the solution of Hierarchy problem, then sparticles must be discovered there [15].

1.3 Supersymmetry

Supersymmetry is a symmetry that relates fermions to bosons that is for every SM fermion there is a scalar partner and for every vector boson of SM there is a fermionic partner [15]. Historically, SUSY was introduced (Haag-Lopuszanski-Sohnius theorem) in a response to Coleman-Mandula theorem that places certain restrictions on the symmetries of S-matrix. SUSY also provides a way to avoid fine tuning problem in Higgs Hierarchy problem. $FCNC$ processes are also possible in $\not{R}_p$ SUSY at tree level, introduced in sects.1 and 1.4.

(a) Coleman-Mandula theorem

Coleman-Mandula theorem is a no-go theorem that provided an impetus to $SUSY$ [16]. This theorem concerns quantum field theories (QFT) with S-matrix that satisfies certain assumptions. Such QFT theories can only have a symmetry of Lie algebra that is always a direct product of the Poincare group and an internal group i.e., no mixing between these two groups is possible. This theorem places constraints on the symmetries of the S-matrix only. Spontaneously broken symmetries are not covered by this theorem.

(b) Haag-Lopuszanski-Sohnius theorem

Haag-Lopuszanski-Sohnius theorem [16] showed that Coleman-Mandula theorem only applies in the case of Lie-algebra. They showed that $SUSY$ is the only additional symmetry allowed by the fact that it is based on Super-Lie-algebra [7].

(c) Wess-Zumino Model

Supersymmetry was revived independently by Yu. A. Golfand and E.P. Likhtman (1971), J. L. Gervais and B. Sakita (1971), D.V. Volkov and V.P. Akulov (1972) and J. Wess and B. Zumino (1974) [16]. Later, A. Salam and J. Strathdee introduced the supergauge transformations [17] to construct SUSY from a simple method.

Consider a mass-less and non-interacting Wess-Zumino model [15]

$$\begin{aligned} S &= \int d^4x(\mathcal{L}_{scalar} + \mathcal{L}_{fermion}), \\ \mathcal{L}_{scalar} &= -\partial^\mu \phi^* \partial_\mu \phi; \mathcal{L}_{fermion} = -i\psi^\dagger \bar{\sigma}^\mu \partial_\mu \psi. \end{aligned}$$

Here, ϕ is a complex wavefunction and ψ is a $Weyl$ spinor given as:

$$\psi = \begin{pmatrix} \varepsilon_\alpha \\ \chi^{\dagger\ \dot{\alpha}} \end{pmatrix},$$

where the dotted and undotted indices help us to distinguish between the complex conjugate self representation and simple self representation of spinor algebra ($SL(2, C)$) [6].

The $SUSY$ generators Q s are then calculated as [15, 19],

$$\begin{aligned} Q_\alpha &= \sqrt{2}\int d^3x J^0_\alpha; \qquad Q^\dagger_{\dot{\alpha}} = \sqrt{2}\int d^3x J^{\dagger 0}_{\dot{\alpha}}, \\ J^\mu_\alpha &= (\sigma^\nu \bar{\sigma}^\mu \psi)_\alpha \partial_\nu \phi^*; \qquad J^{\dagger\mu}_{\dot{\alpha}} = (\psi^\dagger \bar{\sigma}^\mu \sigma^\nu)_{\dot{\alpha}} \partial_\nu \phi. \end{aligned}$$

One can see that the SUSY operators satisfying the following commutation property,

$$\begin{aligned}
\{Q_\alpha, Q^\dagger_{\dot\alpha}\} &= 2\sigma^\mu_{\alpha\dot\alpha} P_\mu;\ (\sigma^0 = 1, \sigma^i = \text{Pauli matrices}) \\
\{Q_\alpha, Q_\beta\} &= \{Q^\dagger_{\dot\alpha}, Q^\dagger_{\dot\beta}\} = 0, \\
[Q_\alpha, P_\mu] &= [Q^\dagger_{\dot\alpha}, P_\mu] = 0.
\end{aligned}$$

$Q\ s$ change the spin of SM particles by 1/2 unit i.e., change SM fermion to a scalar spartner and a vector boson to a fermion spartner,

$$\begin{aligned}
Q|boson &> = |fermion>, \\
Q|fermion &> = |boson> .
\end{aligned}$$

Thus our simple $SUSY$ model contains a $Weyl$ fermion and a scalar boson which form a $SUSY$ multiplet called chiral multiplet. A vector supermultiplet, on the other hand, contains a vector field and a Weyl fermion field. Further, a particle in a supermultiplet contains equal number of bosonic and fermionic degrees of freedom. In SM ϕ and $\phi^\dagger$ provide mass to the fermions (see eq. (1.1.14)). But $\phi^\dagger$ is forbidden in $SUSY$ (see eq. (1.3.1)) [15]. Therefore, we need two independent chiral Higgs supermultiplets $H_u$and H_d.

$$< H_u^0 >= \mathrm{v}_u;\ < H_d^0 >= \mathrm{v}_d.$$

These two parameters are related as

$$\begin{aligned}
\mathrm{v}_u^2 + \mathrm{v}_d^2 &= \mathrm{v}^2 = \frac{2m_Z^2}{g^2 + g'^2} \approx (174\ GeV)^2, \\
\tan\beta &= \frac{\mathrm{v}_u}{\mathrm{v}_d}.
\end{aligned}$$

The 2HDM predicted by SUSY is type II as it satisfies the condition shown by Glashow and Weinberg to prevent tree level contributions to $FCNC$ i.e., H_u doublet couples with up quark only and H_d doublet couples to down type quark only.

(d) Minimal Supersymmetric Standard Model

The Minimal Supersymmetric Standard Model ($MSSM$) [7] was introduced by keeping in view the phenomenological implication of $SUSY$. It contains the minimum number of particles and fields. It is also the minimal extension of SM having $N = 1$ generators. $MSSM$ is constructed as follows [7, 15, 19]:

1. The gauge symmetry group for the theory is selected as the SM gauge group i.e., $SU(3)_C \times SU(2)_L \times U(1)_Y$.
2. Each vector field is assigned to a gauge superfield (VSF) and a fermionic field to a left-chiral superfield (χSF).
3. The Higgs sector is assigned two left chiral scalar superfields having opposite hypercharge.
4. A renormalizable and gauge invariant superpotential is constructed.

(e) Superpotential

The most general interaction $SUSY$ lagrangian, which is also Lorentz invariant and renormalizable, can be written in the following simple form [7, 15, 19]:

$$\mathcal{L}_{int} = -\frac{1}{2}W^{ij}\psi_i\psi_j + c.c.,$$

where W^{ij} is a function of bosonic fields. W^{ij} can also be written in the form

$$W^{ij} = \frac{\delta^2}{\delta\phi_i\delta\phi_j}W,$$

where

$$W = \frac{1}{2}M^{ij}\phi_i\phi_j + \frac{1}{6}y^{ijk}\phi_i\phi_j\phi_k. \tag{1.3.1}$$

W is called the super-potential. It is an analytical function of the complex scalar fields $\phi_i^{,}s$. M^{ij} is fermion mass matrix and y^{ijk} is the Yukawa couplings. In terms of chiral and vector superfields, the superpotential takes the form [15, 19]:

$$\begin{aligned} W &= \varepsilon_{ab}[h^E_{ij}\widehat{H}^a_1\widehat{L}^b_i\widehat{E}^c_j + h^D_{ij}\widehat{H}^a_1\widehat{Q}^b_i\widehat{D}^c_j + h^U_{ij}\widehat{H}^a_2\widehat{Q}^b_i\widehat{U}^c_j - \mu\widehat{H}^a_1\widehat{H}^b_2] + \\ &\quad \varepsilon_{ab}[\frac{1}{2}\lambda_{ijk}\widehat{L}^a_i\widehat{L}^b_j\widehat{E}^c_k + \lambda'_{ijk}\widehat{L}^a_i\widehat{Q}^b_j\widehat{D}^c_k + \mu_j\widehat{H}^a_j\widehat{H}^b_2] + \frac{1}{2}\lambda''_{ijk}\widehat{U}^c_i\widehat{D}^c_j\widehat{D}^c_k. \end{aligned} \tag{1.3.2}$$

Where ε_{ab} are antisymmetric symbols used to raise and lower indices of the spinors.

$$\varepsilon^{12} = \varepsilon_{21} = 1; \varepsilon_{12} = \varepsilon^{21} = -1; \varepsilon_{11} = \varepsilon_{22} = 0.$$

$\widehat{L}$, $\widehat{Q}$, $\widehat{H}_1$ and $\widehat{H}_2$ are left chiral superfield doublets, while $\widehat{E}$, $\widehat{U}$ and $\widehat{D}$ are right chiral superfield singlets as given in table 1.5.3. The term in the first bracket looks familiar with SM lagrangian (see eq. (1.1.14)). These actually describe the $SUSY$ extension of SM. The term $\mu\widehat{H}_1^a\widehat{H}_2^b$ is similar to Higgs mass but it can be made to disappear after rotating the superfields $\widehat{H}_1^a$ and $\widehat{H}_2^b$ [19]. Other terms introduce certain decay processes like proton decay and lepton flavor and number violating processes. To date, no such processes have been discovered. This means that additional symmetries are needed to be introduced for the conservation of these important quantum numbers. R-parity is such a symmetry [15, 7, 19].

(f) R-parity conservation

The second term in superpotential (see eq. (1.3.2)) makes proton decay possible through decay channel ($p^+ \rightarrow \pi^0 e^+, \pi^0 \mu^+$ *etc*, see Fig. (1.5.2)) if the Yukawa couplings λ''_{ijk} and λ'_{ijk} are greater than unity [15].

$$\Gamma(p \rightarrow e^+\pi^0) \sim \frac{m^5_{proton} \sum\limits_{i=2,3} \left|\lambda'_{11i}\lambda''_{11i}\right|^2}{m^4_{\tilde{d}_i}}.$$

Experiments (**Kolar Gold Field - Kolar district** (Kamataka, India); **NUSEX - Mont Blanc** (Alps, France); **FREJUS - Frejus tunnel** (Alps, France); **SOUDAN- Soudan underground mine** (Minnesota, US)) put a maximum limit on proton decay life time is of the order of 10^{34} years [22].

A new symmetry is defined to prevent such terms and to allow other terms in the superpotential that conserve baryon and lepton number. This symmetry must take into account the lepton and baryon number conservation. This symmetry should also take into account the spin quantum number to distinguish between particles and their spartners. R-parity is finally defined as

$$R_p = (-1)^{3(B-L)+2s}.$$

From the table. (1.5.2), it can be seen that R-parity is $+1$ for SM particles and -1 for their spartners. On any vertex, the product of R-parity is $+1$. This leads to many interesting phenomenological consequences:

- Sparticles decay in a cascade manner ending at the lightest sparticle (LSP). This LSP can be a good candidate for dark matter if it does not carry any charge.
- Sparticles are produced only in pairs.

R-parity is an ad-hoc and accidental symmetry. It may be relaxed while assuming that the $\lambda'_{ijk}\lambda''_{lmn}$ product vanishes. Such is the case for R-parity violation.

1.4 R-parity violation

Once R-parity is allowed to be violated keeping the Yukawa couplings involved in proton decay vanishingly small, lepton number and flavor violating decay processes become possible [15, 7, 19].

$$W_{\not{R}_p} = \varepsilon_{ab}[\frac{1}{2}\lambda_{ijk}\widehat{L}_i^a\widehat{L}_j^b\widehat{E}_k^c + \lambda'_{ijk}\widehat{L}_i^a\widehat{Q}_j^b\widehat{D}_k^c] + \frac{1}{2}\lambda''_{ijk}\widehat{U}_i^c\widehat{D}_j^c\widehat{D}_k^c. \tag{1.4.1}$$

Sparticles can now mediate the interaction between the particles of SM [15]. This opens up possibilities for lepton number and flavor violation within MSSM [19]. Flavor changing neutral currents are no longer suppressed as they proceed through tree diagrams within MSSM model [15, 19]. The couplings λ_{ijk}, λ'_{ijk} and λ''_{ijk}'s involved are the parameters of $\not{R}_p$ MSSM. λ_{ijk} is antisymmetric in the first two indices, while λ''_{ijk} is antisymmetric in its last two indices.

$$\lambda_{ijk} = -\lambda_{jik};\ \lambda''_{ijk} = -\lambda''_{ikj}.$$

There are 9 λ_{ijk}, 9 λ''_{ijk} and 27 λ'_{ijk} parameters. $L_iL_jE_k^c$ and $L_iQ_jD_k^c$ operators contribute to pure leptonic and semileptonic decays of hadrons [15]. $\lambda'_{ijk}\lambda_{lmn}$ coupling product make dominant contribution to lepton polarization asymmetry in pure leptonic decays of heavy mesons (charm and beauty), which is zero in SM. Both $\lambda^{'}_{ijk}\lambda_{lmn}$ and $\lambda^{'}_{ijk}\lambda^{'}_{lmn}$ contribute to forward backward asymmetry of leptons measured in the semileptonic decays of hadrons [46].

(a) R-parity violation phenomenology

R-parity violation is rich in phenomenology. Lepton number and flavor violating decay processes are possible in this sector [15, 19].

Neutrinos acquire majorana mass, which constitutes a phenomenology in its own right. The mass is acquired through the tree-level contribution arising from the neutrino-neutralino mixing due to bilinear R-parity violation. Neutrinoless double beta decay, a lepton number violating decay process that may decide the nature of mass for neutrinos (Dirac or Majorana) is also possible within R-parity violation (see Fig. (1.5.3)) [15, 19, 23].

Even the LSP, which is not stable in the presence of $\not{R}_p$ may have a long life-time greater than the age of the universe. This is possible if the Yukawa couplings related to the LSP decay are extremely small. However such a small value of Yukawa couplings makes such processes inaccessible at colliders. An example is [15, 23]

$$\widetilde{\chi}^0 \rightarrow e^+ + 2 \text{ fermions}.$$

The experimental bound on lifetime of the neutralino is [19]

$$\tau(\widetilde{\chi}^0 \rightarrow e^+ + 2 \text{ fermions})/t_0 > 6 \times 10^{10} h (m_{\widetilde{\chi}^0}/100\ GeV)^{1/2} (\widetilde{m}/100\ GeV)^{1/2},$$

where h is the reduced hubble parameter given as [24]

$$H_0 = 100\ h\ km/s/Mpc.$$

Massive neutrinos can have magnetic dipole moments. R-parity violation induces transition magnetic moment i.e., transition between left handed neutrino and right handed anti-neutrino with different flavors [15, 23].

1.5 Searches for R-parity violation at accelerators

Search for R-parity violation has currently been unable to discover any signal. Three facilities **LEP** (Large Electron Positron Collider), **HERA** (Hadron-Elektron-Ringanlage) and **Tevatron** were mainly involved in the search of signals for $SUSY$ and $\not{R}_p$ in the past 18 years [26]. These searches have yet found no evidence for Higgs boson and supersymmetry particles. **LHC** will now take charge of the search for physics beyond SM and including $\not{R}_p$ $MSSM$.

(a) LEP

LEP was a CERN based e^+e^- collider. It started with a total centre of mass (CM) energy of 91 GeV, ultimately reaching 209 GeV at the end of 2000 [25]. It searched for signals of SUSY particles and $\not{R}_p$ MSSM in various decay channels like pair produced Neutralinos ($\widetilde{\chi}^0_{1,2,3,4}$) charginos ($\widetilde{\chi}^{\pm}_{1,2}$), sleptons ($\widetilde{l}_R, \widetilde{v}$) [26]. These pair produced sparticles decay via $\not{R}_p$ Yukawa couplings to SM particles like $\widetilde{v}_i \rightarrow l^+_i l^-_j$. LEP failed to find any signal of SUSY particles and $\not{R}_p$ MSSM during its operation.

(b) HERA

HERA was a Hamburg (DESY) based accelerator. 920 GeV protons were collided with 27 GeV electrons [27]. Various decay processes like the decays of charginos, neutralinos and squarks ($\widetilde{q}$) via $\not{R}_p$ Yukawa couplings were searched [28]. No signal of SUSY and $\not{R}_p$ MSSM was discovered at HERA during its operation.

(c) Tevatron

Tevatron is a Fermi National Accelerator Laboratory (Batavia) based accelerator. Proton and antiproton beam, each of energy 980 GeV are collided to total energy 1.8 TeV [29]. It searched for signals of SUSY and $\not{R}_p$ MSSM in various decay processes like the decays of neutralinos, squarks, sleptons and gluinos ($\widetilde{g}$) [30]. It has yet unable to discover any signal confirming SUSY and $\not{R}_p$ MSSM.

(d) LHC

Large Hadron Collider (LHC), a CERN based accelerator, has been constructed to probe energies at 14 TeV. It has started running since September 2008.

An interesting process is dilepton pair production in $pp \to u_j u_j (d_k d_k) \to l_\alpha^+ l_\alpha^- + X$. An important aspect of this process is that at c.m.s colliding energy beyond M_Z, new physics scenarios like $R\!\!\!/_p$ $MSSM$, extra-dimensions and composite leptons and bosons become accessible [31]. The cross-section of this process is proportional to $(\lambda'_{i33}\lambda_{i\alpha\alpha})$ i.e., involving third generation of quarks. Here, i represents the sneutrino generation index.

Single top quark production is also an interesting prospect for study at LHC with regards to $R\!\!\!/_p$ $MSSM$ [32].

$$pp \to t + g(\widetilde{t} + \widetilde{g}).$$

This process (see Fig. (1.5.5)) violates baryon number within $R\!\!\!/_p$ $MSSM$ and is proportional to λ''. LHC also provides an interesting opportunity to study resonant sneutrino production [33].

$\widetilde{l}$ or a $\widetilde{q}$ can be produced at LHC via λ'_{ijk} $(L_i Q_j D_k)$ or a λ''_{ijk} $(U_i D_j D_k)$ coupling strength respectively. Dominance of λ'_{ijk} $(L_i Q_j D_k)$ leads to Drell-Yan process i.e., decaying to two leptons, while dominance of λ''_{ijk} $(L_i Q_j D_k)$ leads to multijet final states [33]. Following describes resonant sneutrino production via dominant λ'_{ijk} $(L_i Q_j D_k)$ coupling strength:

$$\begin{aligned} pp &\to \widetilde{\nu}_i \to \widetilde{\chi}_1^+ l_i, \\ \widetilde{\chi}_1^+ &\to \widetilde{\chi}_1^0 l^+ \nu;\ \widetilde{\chi}_1^0 \to l_i u_j d_k, \end{aligned}$$

where $\widetilde{\chi}_1^+$ is the chargino and $\widetilde{\chi}_1^0$ is the neutralino.

(e) ILC

International Linear Collider (ILC) is a proposed e^+e^- linear collider of centre of mass energy of 500 GeV. It is expected to be completed in the 2010 decade [34]. The main goal of this future facility is to make a precision measurement of parameters relating to new physics scenarios expected to be discovered at LHC. Supersymmetry and R-parity violation will also be explored at ILC. Top quark pair production ($e^+e^- \to t\ \bar{t}$) [34] is a candidate process to look for signals for R-parity violation (see Fig. (1.5.6)). The cross-section is clearly proportional to $\left|\lambda'_{1i3}\right|^2$, where i is the generation indices for down type squark $\widetilde{d}_i$.

Particles	Mass (GeV)	Mass (GeV)
	Standard Model	Experimental
$W^{\pm}$	80.379	80.399 ± 0.023
Z^0	91.1874	91.1875 ± 0.0021
t	173.2	$173.1 \pm 1.3 \pm 1.1$

Table 1.5.1: Comparison of SM prediction for W, Z and top quark mass with experimental measurements [4].

Particles	Spin	Baryon Number	Lepton Number	(-1)^{2S}	(-1)$^{3(B-L)}$	R_p =(-1)^{2S}(-1)$^{3(B-L)}$
$e, \mu, \tau, \nu_e, \nu_\mu, \nu_\tau$	$\pm\frac{1}{2}$	0	± 1	-1	-1	1
$\gamma, g, W^{\pm}, Z$	1	0	0	1	1	1
u, c, t, d, s, b	$\pm\frac{1}{2}$	1/3	0	-1	-1	1
$\widetilde{e}, \widetilde{\mu}, \widetilde{\tau}, \widetilde{\nu_e}, \widetilde{\nu_\mu}, \widetilde{\nu_\tau}$	0	0	± 1	1	-1	-1
$\widetilde{u}, \widetilde{c}, \widetilde{t}, \widetilde{d}, \widetilde{s}, \widetilde{b}$	0	1/3	0	1	-1	-1
$\widetilde{\gamma}, \widetilde{g}, \widetilde{W^{\pm}}, \widetilde{Z}$	$\frac{1}{2}$	0	0	-1	1	-1

Table 1.5.2: R-parity value for SM and their spartners.

Superfield	SU(3)	$SU(2)_L$	$U(1)_Y$	Multiplet
$\widehat{L}$	1	2	-$\frac{1}{2}$	(e_L, ν_L),$(\widetilde{e}_L, \widetilde{\nu}_L)$
$\widehat{E}^c$	1	1	1	$\widetilde{e}_R, \widetilde{e}^*_R$
$\widehat{Q}$	3	2	1/6	(u_L, d_L),$(\widetilde{u}_L, \widetilde{d}_L)$
$\widehat{U}^c$	3	1	-2/3	$\widetilde{u}_R, \widetilde{u}^*_R$
$\widehat{D}^c$	3	1	1/3	$\widetilde{d}_R, \widetilde{d}^*_R$
$\widehat{H}_1$	1	2	-$\frac{1}{2}$	$(H_1, \widetilde{h}_1)$
$\widehat{H}_2$	1	2	$\frac{1}{2}$	$(H_2, \widetilde{h}_2)$

Table 1.5.3: Chiral superfields(MSSM).

Superfield	SU(3)	$SU(2)_L$	$U(1)_Y$	Multiplet
$\widehat{G}^a$	8	1	0	$g, \widetilde{g}$
$\widehat{W}^i$	1	3	0	$W_i, \widetilde{w_i}$
$\widehat{B}$	1	1	0	$B, \widetilde{b}$

Table 1.5.4: Vector superfields(MSSM).

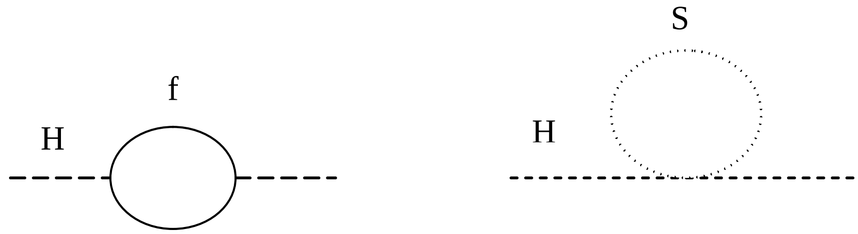


Figure 1.5.1: Virtual fermions and scalar contributing to Higgs mass. Twice contribution from the scalar particles cancel the contributions from one fermion.

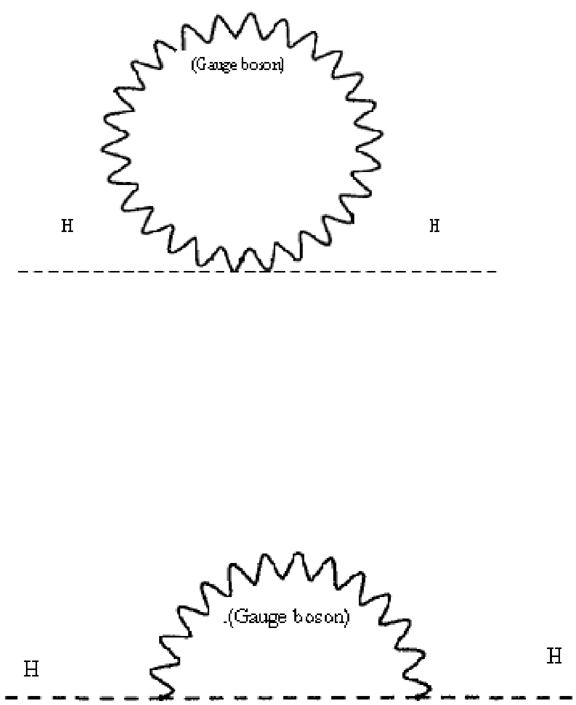


Figure 1.5.2: Gauge boson(A_μ) contribution to Higgs mass.

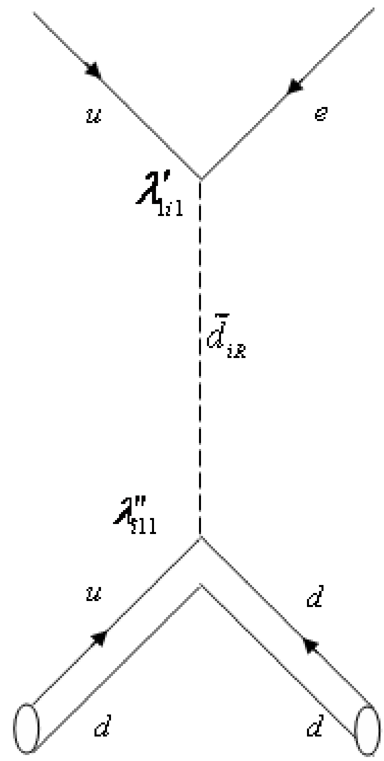


Figure 1.5.3: An example of proton decay $p^+ \to e^+\pi^0$ (Spectator Quark Model).

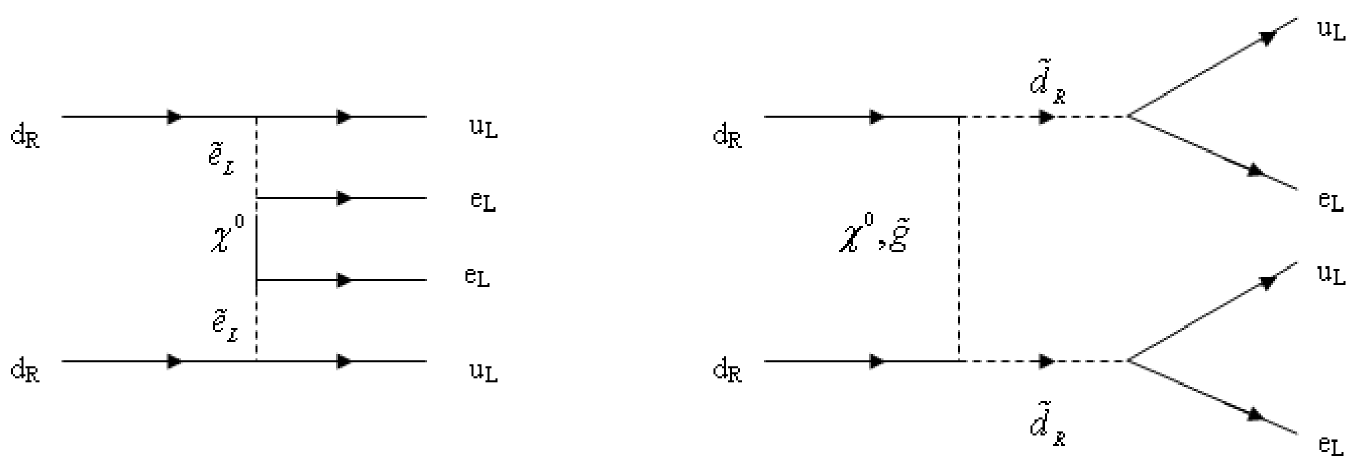


Figure 1.5.4: Neutrinoless double beta decay in $\not{R}_p$ $MSSM$ [19].

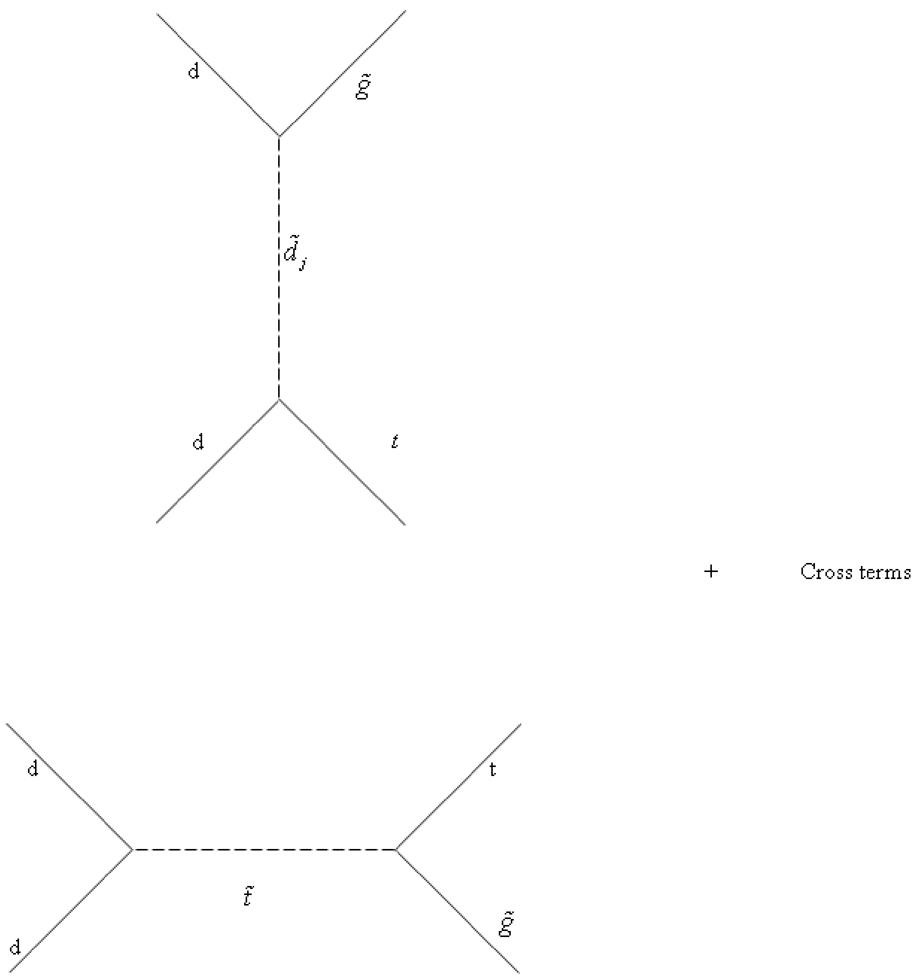


Figure 1.5.5: Feynman diagram of $q_j q_k \rightarrow t\ \tilde{g}$ [32].

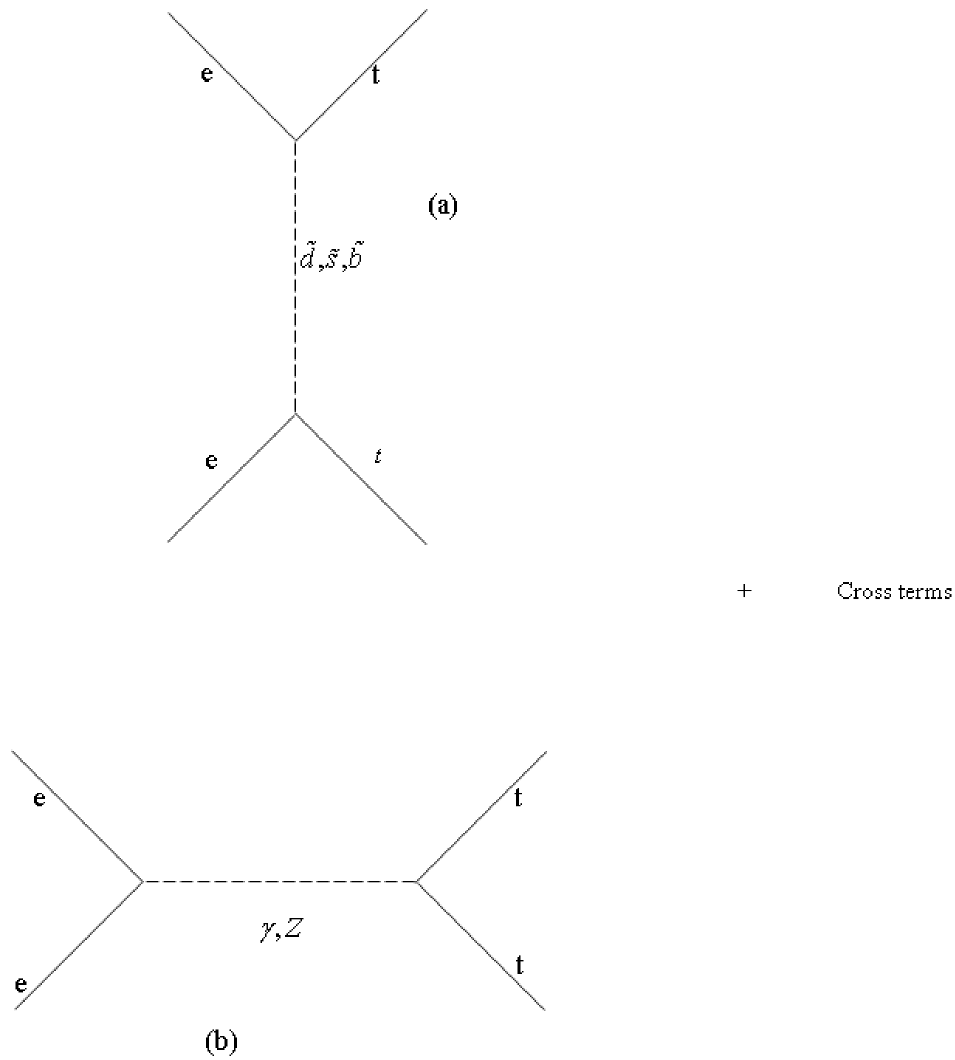


Figure 1.5.6: $e^+e^- \to t\ \bar{t}$ in (a) $R\!\!\!/_p$ $MSSM$. (b) Standard Model [34].

Chapter 2

Semileptonic charged Kaon decays(ΔS=1) in $\not{R}_p$ MSSM

2.1 Introduction

In this chapter we study FCNC semileptonic K meson decays ($K^{\pm} \rightarrow \pi^{\pm} l^{+} l; l = e, \mu$). We derive bounds on R-parity Yukawa couplings. Forward backward asymmetry and lepton polarization asymmetry related to these decay processes is also studied.

K meson decays are one of the most important phenomenological tools to study different aspects of SM like CP-violation and elements of CKM matrix [35]. They also provide examples where one may study physics beyond SM [36]. Semileptonic decays of mesons are one such example. These decays are forbidden in SM at tree level [6]. However, higher order diagrams like box and penguin do contribute to these processes in the SM [6]. Leptonic K meson decays have been studied especially for the determination of CKM matrix elements and CP-violation [37].

Flavor changing neutral current K decays like ($K^0 \rightarrow \pi^0 l^+ l^-$, $K^{\pm} \rightarrow \pi^{\pm} \nu\nu$) are also clean probe of important aspects of SM like flavor structure [37, 38]. MSSM contributes significantly to these decays [38]. These decays have been calculated in SM by several authors [39], and have also been studied by others within 2HDM (2 Higgs Doublet Model) and extra dimensions [40]. These decays have been probed in CERN experiment NA48 [41] and a new experimental facility NA48/3 has also been proposed recently to study such decays [41].

2.2 Semileptonic charged Kaon decays

In this chapter we will study SM and $\not{R}_p$ contribution to semileptonic charged kaon decays ($K^{\pm} \rightarrow \pi^{\pm} l^{+} l^{-}$). The $\not{R}_p$ $MSSM$ interaction lagrangian in terms of component fields [18, 19]:

$$\begin{aligned} L_{\not{R}_p} &= \lambda_{ijk}[\widetilde{\nu}_{i_L}\bar{e}_{k_R}e_{j_L} + \widetilde{e}_{j_L}\bar{e}_{k_R}\nu_{i_L} + \widetilde{e}^*_{k_R}\overline{(\nu_{i_L})^c}\bar{e}_{j_L}] + \\ & \lambda'_{ijk}\left[\begin{array}{c} \widetilde{\nu}_{i_L}\bar{d}_{k_R}d_{j_L} + \widetilde{d}_{j_L}\bar{d}_{k_R}\nu_{i_L} + \widetilde{d}^*_{k_R}\overline{(\nu_{i_L})^c}d_{j_L} \\ -\nu^{\dagger}_{jp}(\widetilde{e}_{i_L}\bar{d}_{k_R}u_{p_L} + \widetilde{u}_{j_L}\bar{d}_{k_R}e_{i_L} + \widetilde{d}^*{}_{k_R}\overline{(e_{i_L})^c}u_{p_L}) \end{array}\right] + h.c, \end{aligned} \tag{2.2.1}$$

where d is three generation down quark; while u corresponds to first two generation up quark.

The corresponding effective Lagrangian in the quark mass basis is as follows [18, 19]:

$$L^{eff}_{\not{R}_p}\left(\overline{s} \longrightarrow \overline{d} + l_\beta + \overline{l}_\beta\right) = \frac{G_F}{\sqrt{2}}\left[\begin{array}{c} A^{ds}_{\alpha\beta}\left(\overline{l_\beta}\gamma^\mu P_L l_\beta\right)\left(\overline{q}_s\gamma_\mu P_R q_d\right) \\ -B^{ds}_{\alpha\beta}\left(\overline{l_\beta}P_R l_\beta\right)\left(\overline{q}_s P_L q_d\right) \\ -C^{ds}_{\alpha\beta}\left(\overline{l_\beta}P_L l_\beta\right)\left(\overline{q}_s P_R q_d\right) \end{array}\right]. \tag{2.2.2}$$

$\beta = e,\ \mu$. The dimensional coupling constants $A^{ds}_{\alpha\beta},\ B^{ds}_{\alpha\beta},\ C^{ds}_{\alpha\beta}$, depend on the species of quark and charged leptons:

$$A^{ds}_{\alpha\beta} = \frac{\sqrt{2}}{G_F}\sum_{m,n,i=1}^{3}\frac{V^{\dagger}_{ni}V_{im}}{2m^2_{\widetilde{u^c_i}}}\lambda'_{\beta n1}\lambda'^*_{\alpha m2} \tag{2.2.3}$$

$$B^{ds}_{\alpha\beta} = \frac{\sqrt{2}}{G_F}\sum_{i=1}^{3}\frac{2}{m^2_{\widetilde{\nu}_{Li}}}\lambda^*_{i\alpha\beta}\lambda'_{isd}, \tag{2.2.4}$$

$$C^{ds}_{\alpha\beta} = \frac{\sqrt{2}}{G_F}\sum_{i=1}^{3}\frac{2}{m^2_{\widetilde{\nu}_{Li}}}\lambda_{i\beta\alpha}\lambda'^*_{isd}. \tag{2.2.5}$$

The effective Hamiltonian for $K \rightarrow \pi l^{+} l^{-}$is given in SM as [42, 43]

$$H_{eff}(\Delta S = 1) = -\frac{G_F}{\sqrt{2}}\sum_{i=u,c,t}V^*_{is}V_{id}[\frac{\alpha}{2\pi}\frac{Y(x_i)}{(\sin\theta_W)^2}](\bar{s}d)_{V-A}(\bar{\mu}\mu)_{V-A}] + H.c., \tag{2.2.6}$$

where $x_i = (\frac{m_i}{M_W})^2$ and

$$\begin{aligned}
Y(x) &= Y_0(x) + \frac{\alpha_s}{4\pi} Y_1(x) \\
Y_0(x) &= \frac{x}{8}(\frac{4-x}{1-x} + \frac{3x}{(1-x)^2}\ln[x]) \\
Y_1(x) &= \frac{x^3+2x}{(x-1)^2} Li_2(1-x) + \frac{x^4 - x^3 + 14x^2 - 2x}{2(x-1)^3}\log^2 x \\
&\quad - \frac{x^4 + x^3 + 10x^2 - 4x}{(x-1)^3}\log x + \frac{4x^3 + 16x^2 + 4x}{3(x-1)^2} \\
&\quad + [\frac{2x^2 - 4x}{x-1} - \frac{x^3 - 7x^2}{(x-1)^2} - \frac{6x^2}{(x-1)^3}\log x]\log(\frac{\mu^2}{M_W^2}),
\end{aligned} \tag{2.2.7}$$

where α_s is the strong coupling constant and μ is the renormalization scale, at which the top quark mass is renormalized. $Y_1(x)$ represents the NLO gluonic correction to box and penguin diagram [49].

The most general invariant amplitude for the decay ($K^{\pm} \to \pi^{\pm} l^+ l^-$) can be written as [46]

$$M = F_S \bar{l} l + F_V (p_K)_\mu \bar{l}\gamma^\mu l + F_A (p_K)_\mu \bar{l}\gamma^\mu\gamma^5 l + F_p \bar{l}\gamma^5 l, \tag{2.2.8}$$

where F_A, F_V, F_P and F_S are coefficients of axial-vector, vector, pseudo-scalar and scalar interaction terms respectively. Following appendix A and substituting eqs. (2.2.1-2.2.6) in eq. (2.2.8), we use form factors $f_+(s)$ as unity and $f_-(s)$ neglected at the lowest orders [44, 46]. The coefficients are then given by comparing eqs. (A.1) with eq. (2.2.2) and then using eq. (A.6).

$$\begin{aligned}
F_V &= f^+(s) C_{VV} = -\frac{G_F}{\sqrt{2}} \sum_{i=u,c,t} \frac{\alpha}{2\pi} V_{is}^* V_{id} \frac{Y(x_i)}{(\sin\theta_W)^2} + \sum_{m,n,i=1}^{3} \frac{V_{ni}^\dagger V_{im}}{2m^2_{\widetilde{u_i^c}}} \lambda'_{\beta nk}\lambda'^*_{\alpha mp}, \qquad (2.2.9) \\
F_A &= f^+(s) C_{VA} = \frac{G_F}{\sqrt{2}} \sum_{i=u,c,t} \frac{\alpha}{2\pi} V_{is}^* V_{id} \frac{Y(x_i)}{(\sin\theta_W)^2} - \sum_{i,m,n=1}^{3} \frac{V_{ni}^\dagger V_{im}}{2m^2_{\widetilde{u_i^c}}} \lambda'_{\beta nk}\lambda'^*_{\alpha mp}, \\
F_S &= \frac{m_K^2}{m_s} f_o(s)\ C_{SS} = \sum_{i=1}^{3} \frac{m_K^2}{m_s} \frac{(\lambda^*_{i\alpha\beta}\lambda'_{ipk} + \lambda_{i\beta\alpha}\lambda'^*_{ikp})}{2m^2_{\widetilde{\nu}_{Li}}}, \\
F_P &= -i\frac{m_K^2}{m_s} f_o(s)\ C_{AA} = -i\sum_{i=1}^{3} \frac{m_K^2}{m_s} \frac{(\lambda^*_{i\alpha\beta}\lambda'_{ipk} - \lambda_{i\beta\alpha}\lambda'^*_{ikp})}{2m^2_{\widetilde{\nu}_{Li}}},
\end{aligned}$$

The Forward-Backward Asymmetry (A_{FB}) is defined as the asymmetric angular distribution of dilepton pair with respect to the initial meson direction of momentum in the dilepton rest frame. The A_{FB} is given as:

$$A_{FB}(s) = \frac{\int_0^1 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta}}{\int_0^1 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta} + \int_{-1}^0 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta}}, \tag{2.2.10}$$

where Γ is the decay rate of $K^\pm \to \pi^\pm l^+ l^-$. In the dilepton rest frame [46](see eq. (B.6))

$$A_{FB}(s) = \frac{1}{128\pi^3 m_K^3 (\frac{d\Gamma}{ds})} m_l \beta(m_l, s)^2 \lambda(s)\,\mathrm{Re}(F_S F_V^*), \tag{2.2.11}$$

where

$$\frac{d\Gamma}{ds} = \frac{1}{256\pi^3 m_K^3} \beta(m_l, s) \lambda^{\frac{1}{2}}(s) R(s), \tag{2.2.12}$$

$$\beta(m_l, s) = (1 - \frac{4m_l^2}{s})^{\frac{1}{2}},$$

$$\lambda(s) = m_K^4 + m_\pi^4 + s^2 - 2sm_\pi^2 - 2sm_K^2 - 2m_\pi^2 m_K^2,$$

where

$$\begin{aligned} R(s) \;=\; & |F_S|^2\, 2s\beta(m_l, s)^2 + |F_P|^2\, 2s + |F_V|^2 \frac{1}{3}\lambda(s)(1 + \frac{2m_l^2}{s}) + |F_A|^2 [\frac{1}{3}\lambda(s)(1 + \frac{2m_l^2}{s}) + 8m_K^2 m_l^2] \\ & + \mathrm{Re}(F_P F_A^*) 4m_l(m_K^2 - m_\pi^2 + s), \end{aligned}$$

and s (invariant mass squared of the dilepton system in its rest frame) is bounded as:

$$4(m_l)^2 \le s \le (m_K - m_\pi)^2.$$

In the SM, A_{FB} vanishes in $(K \to \pi l\bar{l})$ as the decay amplitude involves no scalar term [46]. Therefore, any measurement of non-zero A_{FB} in this decay process will give a signal of new physics. The possibility of non-zero A_{FB} has been reported by [47] in case of large tanβ for MSSM. Some authors have also explored the possibility of non-zero A_{FB} in

$R\!\!\!/_p$ framework [45, 48].

We discuss similar possibility in $(K \to \pi l\bar{l})$ within $R\!\!\!/_p$ framework. We show that A_{FB} arises naturally by sneutrino and squark exchange terms. The branching fraction of the decay process $(K \to \pi l\bar{l})$ is given as

$$BR(K^{\pm} \to \pi^{\pm} l\bar{l}) = \int_{4(m_l)^2}^{(m_K - m_\pi)^2} \frac{d\Gamma}{ds} ds\ \tau_K. \tag{2.2.13}$$

where τ_K is the life-time of charged K-meson. Also the inclusive branching fraction can be calculated [19]

$$\Gamma\left[K^{\pm} \to \pi^{\pm} l^+ l^-\right] = \frac{m_K^5}{192\pi^3} G_F^2 \{\mid A^{ds}_{\alpha\beta} \mid^2 + \frac{1}{4}(\mid B^{ds}_{\alpha\beta} \mid^2 + \mid C^{ds}_{\alpha\beta} \mid^2)\}. \tag{2.2.14}$$

2.3 Results and discussion

Tables. (2.3.1 and 2.3.2) display bounds on Yukawa couplings calculated using eq. (2.2.14) and eq. (2.2.12). They also compares these bounds with the previous bounds given in tables (2.3.1 and 2.3.2) [18, 19]. Single coupling dominance has been assumed in the calculation of these bounds.

Considering the bounds calculated from inclusive decay mode $(K^+ \longrightarrow \pi^+ e^+ e^-)$, common limits 1.02×10^{-4} (1σ) and 1.04×10^{-4} (2σ) on the magnitude of three coupling products $\left|\lambda'_{1i1}\lambda'^*_{1i2}\right|$ $(i = 1, 2, 3)$ are obtained. These limits are comparable to those obtained previously for pure leptonic K meson decays [18, 19]. Similarly, common limit 5.10×10^{-5} (1σ) and 5.20×10^{-5} (2σ) on the magnitude of four coupling products $\left|\lambda_{i11}\lambda'^*_{i21}\right|$ and $\left|\lambda^*_{i11}\lambda'_{i12}\right|$ respectively. These limits are much larger $\sim O(10^3)$ to those obtained previously for pure leptonic K meson decays [18, 19].

For $K^+ \longrightarrow \pi^+ \mu^+ \mu^-$, three coupling products of $\left|\lambda'^*_{2i1}\lambda'_{2i2}\right|$ have the common limit 5.73×10^{-5} (1σ), 6.13×10^{-5} (2σ). These limits are larger than the previous bounds $\sim O(10)$. Further, common limits 2.86×10^{-5} (1σ), 3.07×10^{-5} (2σ) on four coupling products $\left|\lambda_{i22}\lambda'^*_{i12}\right|$ and $\left|\lambda^*_{i22}\lambda'_{i21}\right|$ $(i = 1, 3)$, which are larger than the previous bounds $\sim O(10^2)$.

For $K^+ \longrightarrow \pi^+ \mu^+ e^-$, we obtain comparable limits 9.83×10^{-7}on $\left|\lambda'_{2i1}\lambda'^*_{1i2}\right|$, to the previous bounds. We also obtain limits 4.91×10^{-7}on $\left|\lambda_{i21}\lambda'^*_{i12}\right|$, which are larger than the previous bounds $\sim O(10^2)$.

We have also studied these bounds from an exclusive mode along with inclusive mode shown in Figs. (2.3.3-2.3.3) by using eqs. (2.2.6-2.2.11). The results have been shown in Figs. (2.3.2-2.3.3). Under the assumption that only one product combination is non zero, we obtain the limit 4.77×10^{-5} (1σ) and 4.88×10^{-5} (2σ) on the magnitude of coupling product $\left|\lambda'_{131}\lambda'^{*}_{132}\right|$. These are comparable to the previous bounds. We also obtain limit 2.16×10^{-6} (1σ) and 2.21×10^{-6} (2σ) on the magnitude of coupling products $\left|\lambda_{311}\lambda'^{*}_{321}\right|$ and $\left|\lambda^{*}_{311}\lambda'_{312}\right|$ for the decay $K^+ \longrightarrow \pi^+ e^+ e^-$. These limits are smaller from the one calculated from the inclusive mode [20] $\sim O(10)$ and larger than the previous bounds $\sim O(10^2)$.

For $K^+ \longrightarrow \pi^+\mu^+\mu^-$, where the coupling product of $\left|\lambda'^{*}_{231}\lambda'_{232}\right|$ have the limit 9.49×10^{-5} (1σ) and 1.17×10^{-4} (2σ). These are comparable to those bounds obtained by assuming inclusive decay mode. We further obtain limit 2.25×10^{-5} (1σ) and 2.41×10^{-5} (1σ), on coupling products $\left|\lambda_{322}\lambda'^{*}_{312}\right|$ and $\left|\lambda^{*}_{322}\lambda'_{321}\right|$. These limits are almost identical to the one calculated from the inclusive decay mode [20].

Eqs. (2.2.11,2.2.12) have been used to calculate A_{FB}. Vector and axial coefficient $F_V, F_A(\lambda'\lambda')$ are induced by squark exchange while scalar coefficient $F_S(\lambda'\lambda)$ are generated by the sneutrino exchange. We have assumed value of coefficients $F_S,\ F_V,\ F_P,\ F_A \ \sim\ O(10^{-9})$, which have been estimated from tables. (2.3.1 and 2.3.2) i.e., bounds on Yukawa Coupling products. The A_{FB} calculated is of $\sim O(10^{-3})$ for $K^+ \longrightarrow \pi^+ e^+ e^-$ and $\sim\ O(10^{-1})$ for $K^+ \longrightarrow \pi^+\mu^+\mu^-$ (see Figs. (2.3.4-2.3.5)) respectively as calculated in [20].

We have compared bounds on the products of Yukawa couplings for exchange of squarks and sneutrinos in the process $K^+ \longrightarrow \pi^+ l\bar{l}$ in the mass scale $\sim O(100\ GeV)$. The estimated A_{FB} is vanishingly small for $K^+ \longrightarrow \pi^+ e^+ e^-$ due to smallness of electron mass but is large as $\sim O(10^{-1})$ for $K^+ \longrightarrow \pi^+\mu^+\mu^-$ as calculated in both exclusive and inclusive decay mode. This we believe is a measurable effect in future experiments searching for $K^{\pm} \rightarrow \pi^{\pm}\mu^+\mu^-$.

Process	Combination constrained Q_i	Bounds (Inclusive Mode)	Combination constrained Q_i	Bounds (Exclusive Mode)	Previous Bounds
$d\bar{s} \to e^+e^-$	$\lambda'_{111}\lambda'^*_{112}$	$1.02 \times 10^{-4}(1\sigma)$	$\lambda'_{131}\lambda'^*_{132}$	$4.77 \times 10^{-5}(1\sigma)$	8.1×10^{-5}
	$\lambda'_{121}\lambda'^*_{122}$	$9.74 \times 10^{-5}(-1\sigma)$		$4.56 \times 10^{-5}(-1\sigma)$	
	$\lambda'_{131}\lambda'^*_{132}$	$1.04 \times 10^{-4}(2\sigma)$		$4.88 \times 10^{-5}(2\sigma)$	
		$9.51 \times 10^{-5}(-2\sigma)$		$4.46 \times 10^{-5}(-2\sigma)$	
$d\bar{s} \to \mu^+\mu^-$	$\lambda'_{211}\lambda'^*_{212}$	$5.73 \times 10^{-5}(1\sigma)$	$\lambda'_{231}\lambda'^*_{232}$	$9.49 \times 10^{-5}(1\sigma)$	7.8×10^{-6}
	$\lambda'_{221}\lambda'^*_{222}$	$4.8 \times 10^{-5}(-1\sigma)$		$7.97 \times 10^{-5}(-1\sigma)$	
	$\lambda'_{231}\lambda'^*_{232}$	$6.13 \times 10^{-5}(2\sigma)$		$1.17 \times 10^{-4}(2\sigma)$	
		$4.23 \times 10^{-5}(-2\sigma)$		$7.09 \times 10^{-5}(-2\sigma)$	
$d\bar{s} \to e^-\mu^+$	$\lambda'_{211}\lambda'^*_{112}$	9.83×10^{-7}			3×10^{-7}
	$\lambda'_{221}\lambda'^*_{122}$				
	$\lambda'_{231}\lambda'^*_{132}$				
$d\bar{s} \to e^+\mu^-$	$\lambda'_{111}\lambda'^*_{212}$	4.24×10^{-6}			3×10^{-7}
	$\lambda'_{121}\lambda'^*_{222}$				
	$\lambda'_{131}\lambda'^*_{232}$				

Table 2.3.1: Limits on the magnitudes of combination constraints on Yukawa couplings involving various flavors of quarks and leptons with exchange of supquarks; $Q_i = 1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.

Process	Combination constrained L_i	Bounds (Inclusive Mode)	Combination constrained L_i	Bounds (Exclusive Mode)	Previous Bounds
$d\bar{s} \to e^+e^-$	$\lambda_{211}\ \lambda^{'*}_{212}$	$5.10 \times 10^{-5}(1\sigma)$	$\lambda^*_{311}\lambda^{'}_{321}$	$2.16 \times 10^{-6}(1\sigma)$	1.0×10^{-8}
	$\lambda_{311}\ \lambda^{'*}_{312}$	$4.87 \times 10^{-5}(-1\sigma)$	$\lambda^*_{311}\lambda^{'}_{321}$	$2.07 \times 10^{-6}(-1\sigma)$	
	$\lambda^*_{211}\ \lambda^{'}_{221}$	$5.20 \times 10^{-5}(2\sigma)$		$2.21 \times 10^{-6}(2\sigma)$	
	$\lambda^*_{311}\ \lambda^{'}_{321}$	$4.76 \times 10^{-5}(-2\sigma)$		$2.08 \times 10^{-6}(-2\sigma)$	
$d\bar{s} \to \mu^+\mu^-$	$\lambda_{122}\ \lambda^{'*}_{112}$	$2.86 \times 10^{-5}(1\sigma)$	$\lambda_{322}\lambda^{'*}_{312}$	$2.25 \times 10^{-5}(1\sigma)$	2.2×10^{-7}
	$\lambda_{322}\ \lambda^{'*}_{312}$	$2.4 \times 10^{-5}(-1\sigma)$	$\lambda^*_{322}\lambda^{'}_{321}$	$1.89 \times 10^{-5}(-1\sigma)$	
	$\lambda^*_{122}\ \lambda^{'}_{121}$	$3.07 \times 10^{-5}(2\sigma)$		$2.41 \times 10^{-5}(2\sigma)$	
	$\lambda^*_{322}\ \lambda^{'}_{321}$	$2.12 \times 10^{-5}(-2\sigma)$		$1.68 \times 10^{-5}(-2\sigma)$	
$d\bar{s} \to e^-\mu^+$	$\lambda_{121}\ \lambda^{'*}_{112}$	4.91×10^{-7}			6×10^{-9}
	$\lambda_{321}\ \lambda^{'*}_{312}$				
	$\lambda^*_{212}\ \lambda^{'}_{221}$				
	$\lambda^*_{312}\ \lambda^{'}_{321}$				
$d\bar{s} \to e^+\mu^-$	$\lambda_{112}\ \lambda^{'*}_{112}$	2.12×10^{-6}			6×10^{-9}
	$\lambda_{312}\ \lambda^{'*}_{312}$				
	$\lambda^*_{221}\ \lambda^{'}_{221}$				
	$\lambda^*_{321}\ \lambda^{'}_{321}$				

Table 2.3.2: Limits on the magnitudes of combination constraints on Yukawa couplings involving various flavors of quarks and leptons with exchange of sneutrinos; $L_i = 1/(m_{\tilde{v}_{i_L}}/100GeV)^2$.

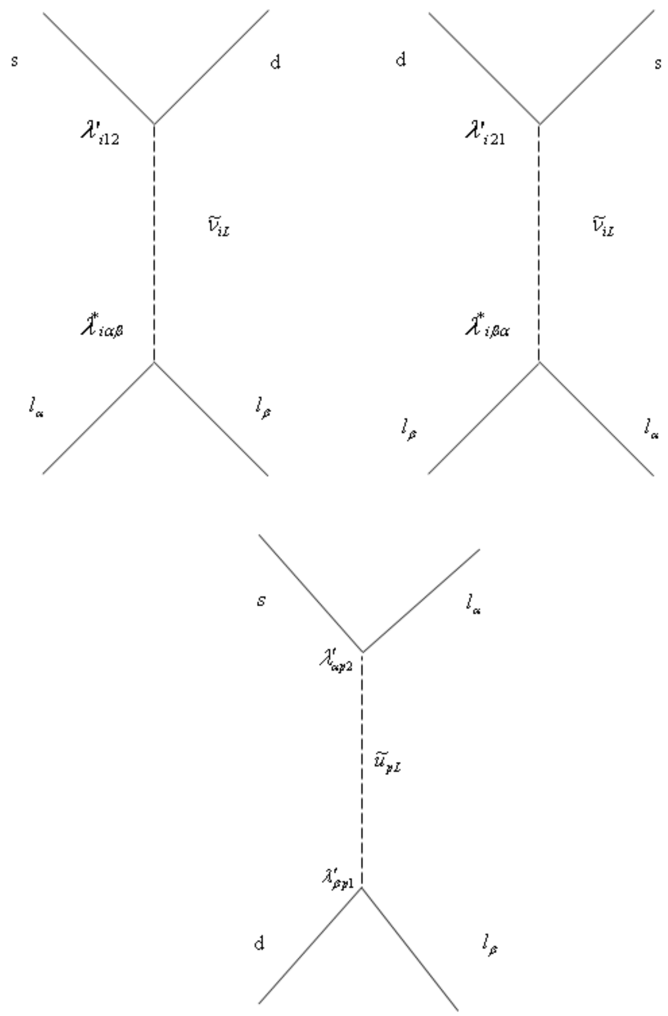
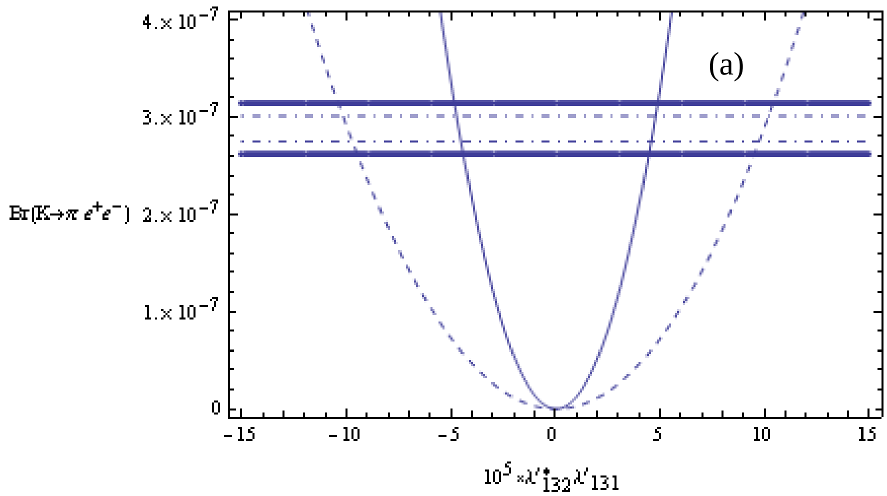


Figure 2.3.1: Tree-Level diagrams contributing to charged Kaon decays in the spectator quark model.
$(K^{\pm} \rightarrow \pi^{\pm} l^{+} l^{-} (\bar{s} \rightarrow \bar{d} l^{+} l^{-}); l = e, \mu)$



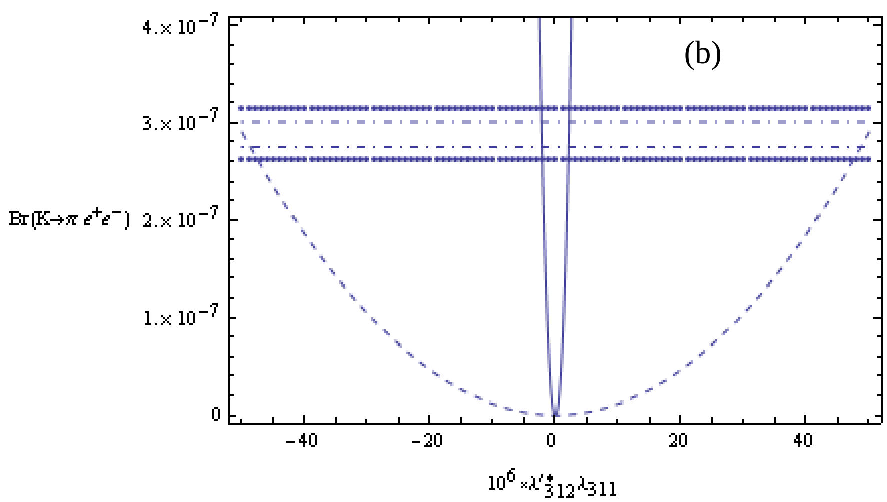


Figure 2.3.2: Variation of Branching fraction of $K^{\pm} \rightarrow \pi^{\pm}e^{+}e^{-}$ w.r.t. $\lambda'^{*}_{132}\lambda'_{131}(a)$ and $\lambda'^{*}_{312}\lambda_{311}(b)$. The dotted curve represents the variation (inclusive mode). The horizontal lines represents the observed branching fraction with 1σ (dashed) and 2σ(double dashed) error. $\lambda'^{*}_{312}\lambda_{311}$ is normalized to $1/(m_{\tilde{v}_{3_L}}/100GeV)^2$. $\lambda'^{*}_{132}\lambda'_{131}$ is normalized to $1/(m_{\tilde{u}_{3_L}}/100GeV)^2$.

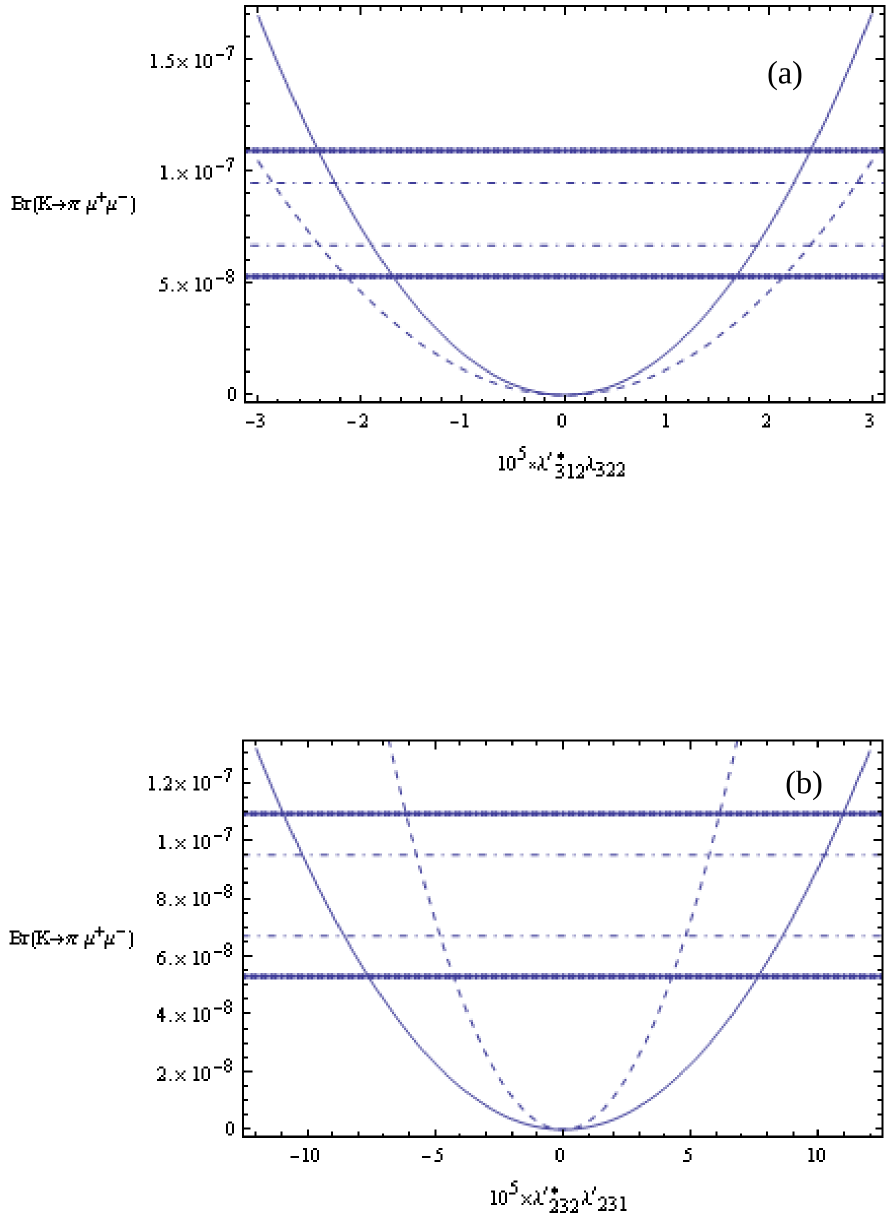


Figure 2.3.3: Variation of Branching fraction of $K^{\pm} \rightarrow \pi^{\pm}\mu^{+}\mu^{-}$ w.r.t. $\lambda'^{*}_{312}\lambda_{322}(a)$ and $\lambda'^{*}_{232}\lambda'_{231}(b)$. The dotted curve represents the variation (inclusive mode). The horizontal lines represents the observed branching fraction with 1σ (dashed) and 2σ(double dashed) error. $\lambda'^{*}_{312}\lambda_{322}$ is normalized to $1/(m_{\tilde{v}_{3_L}}/100GeV)^2$.$\lambda'^{*}_{232}\lambda'_{231}$ is normalized to $1/(m_{\tilde{u}_{3_L}}/100GeV)^2$.

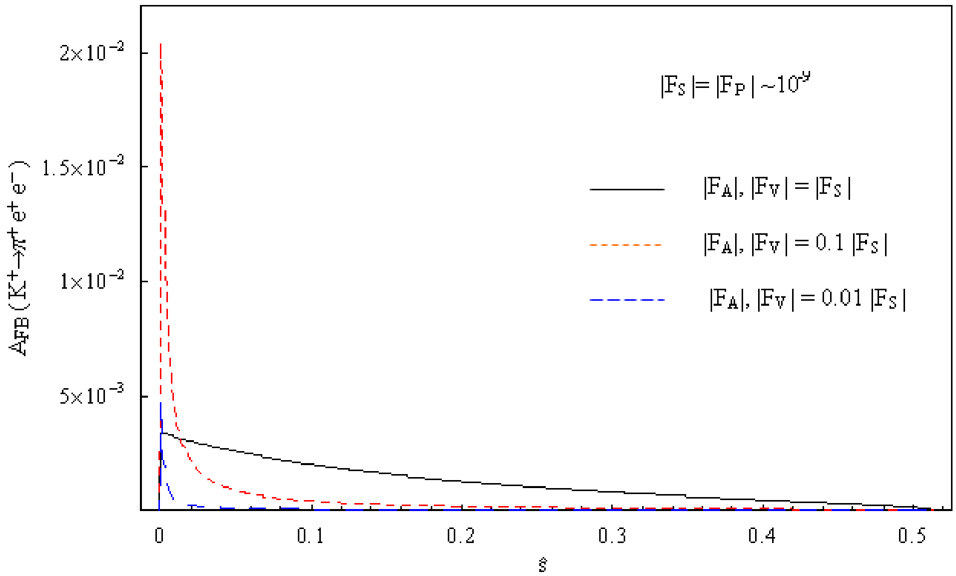


Figure 2.3.4: Variation of forward-backward asymmetry in decay process $(K^+ \to \pi^+ e^+ e^-)$, varying F_V and F_A with constant F_S and F_P, as a function of $\hat{s} = \frac{s}{m_K^2}$.

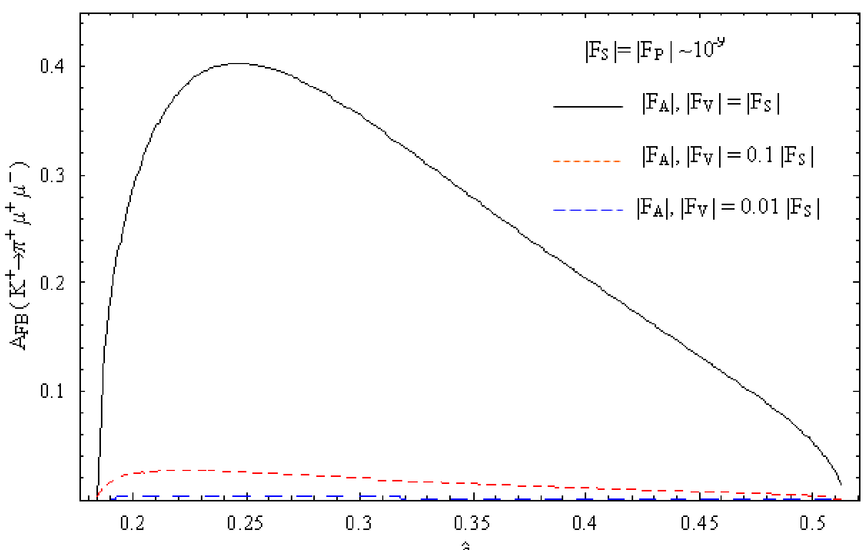


Figure 2.3.5: Variation of forward-backward asymmetry in decay process $(K^+ \to \pi^+ \mu^+ \mu^-)$, varying F_V and F_A with constant F_S and F_P, as a function of $\hat{s} = \frac{s}{m_K^2}$.

Chapter 3

Pure Leptonic decays of B and K mesons in $\not{R}_p$ MSSM

3.1 Introduction

SM has been successful in making accurate and precise predictions about nearly all observations (except for neutrino oscillations) in particle physics. It is still considered as a low energy effective limit of a more general theory.

FCNC processes are suppressed in SM as they are not allowed at tree level in SM and proceed through higher order diagrams (box and penguin). So observables relating to FCNC B and K meson decays receive interesting contributions from physics beyond SM. These decay processes have also played a key role in other areas of SM like measurements of Cabibbo-Kobayashi-Maskawa (CKM) unitary angles [21], probing CP-violation [50]. Belle and BaBar experiments have made significant contribution to this promising and potentially important area of particle physics. Super B-factories, which are a new generation asymmetric high energy e^+e^- colliders of luminosity 10^{36} $cm^{-2}s^{-1}$ [51], can prove to be a new probe for new physics in the flavor sector of the SM. These factories will search for the effects of physics beyond the SM. Study of effects of CP-violating asymmetries in the decays of B and D mesons, and in very rare heavy quark and tau lepton decays will be a part of the search of B factories [51]. Various studies have explored the consequences of the contribution of NP, like 2 Higgs doublet Model ($2HDM$), $MSSM$ and extra dimensions, etc, to B decays [52]. $B\overline{B}$ mixing and leptonic B-decays are one of the many interesting decay processes present in B-physics [53]. Leptonic B decays may provide a clean signal of physics beyond SM. Branching ratio in these decay processes directly provides bounds on NP parameters. Observables like Polarization asymmetry may also show sensitivity to NP parameters. Various studies have considered NP contribution to these observables in different models like $2HDM$ and extra dimensions and confirmed the sensitivity of these observables to the parameters of NP [54, 55]. Pure leptonic K meson decays have been used in the study of $\Delta S = \Delta Q$ rule [56]. These decays have also been calculated in SM [43].

In this chapter, we shall discuss the theoretical predictions of branching ratio and polarization asymmetry of FCNC ($M(p\bar{q}) \rightarrow l\bar{l};\ l = e, \mu$) on the basis of $\not{R}_p$ in the light of current experimental data [73]. We explore this possibility in the framework of $\not{R}_p$ $MSSM$. Here, these interactions arise naturally as sneutrino and squark exchange terms.

3.2 Pure Leptonic meson decays

The effective Hamiltonian for the ($p \rightarrow q\ l_\beta^+\ l_\beta^-;\ \beta = e, \mu$) is given in terms of four Fermi interactions using notation [63, 64].

$$H_{eff} = \sum_i C_{ii}(\bar{q}\Gamma_i b)(\bar{l}_\beta \Gamma_i l_\beta) + \sum_i C'_{ii}(\bar{q}\Gamma_i b)(\bar{l}_\beta \Gamma'_i l), \tag{3.2.1}$$

which involves interactions having coefficients C_{ii} arising from SM and new physics , where Γ_i is one of the Dirac matrices $\{I, \gamma_5, \gamma^\mu, \gamma^\mu\gamma_5, \sigma_{\mu\nu}\}, \Gamma'_i = \Gamma_i\ \gamma_5$ and $i = \{S, P, V, A, T\}$.

The effective Hamiltonian for $B_q^\pm \rightarrow l^+ l^-$ is given in SM as [43]

$$H_{eff} = -\frac{G_F}{\sqrt{2}} \sum_{i=u,c,t} V_{ip}^* V_{iq} [\frac{\alpha}{2\pi} \frac{Y(x_i)}{(\sin\theta_W)^2}](\bar{p}q)_{V-A}(\bar{l}l)_{V-A}] + H.c. \tag{3.2.2}$$

For ($B \rightarrow l_\beta^+ l_\beta^-$), the matrix element $< 0|\bar{q}\Gamma_i b|B >$ vanish where Γ_i =$\{I, \gamma_\mu, \sigma_{\mu\nu}\}$. The matrix element of $(B \to l_\beta^+ l_\beta^-)$ is then given as [61]

$$M = C_{PP}\ (\bar{p}\gamma^5 q)(\overline{l_\beta}\gamma^5 l_\beta) + C_{PS}\ (\bar{p}\gamma^5 q)(\overline{l_\beta} l_\beta) + C_{AA}\ (\bar{p}\gamma^\mu\gamma^5 q)(\overline{l_\beta}\gamma_\mu\gamma^5 l_\beta), \tag{3.2.3}$$

where $C_{PP},\ C_{PS},\ C_{AA}$ are coefficients of pseudoscalar, pseudoscalar scalar and axial interaction terms as given in [64] These couplings will be calculated in $R\!\!\!/_p$ MSSM in this paper. Using PCAC we can calculate branching fraction and A_{LP} (see eq. (C.7)) of the decay process $(M \to l_\beta^+ l_\beta^-)$ [61, 62]

$$\begin{aligned} Br\left[M \to l_\beta^+ l_\beta^-\right] &= \frac{1}{8\pi}(f_{B_q})^2 (m_M)^3 \tau_B \sqrt{1 - \frac{4(m_{l_\beta})^2}{(m_M)^2}} \\ & [\quad | \quad 2\frac{m_{l_\beta}}{m_M} C_{AA} - \frac{m_M}{m_p + m_q} C_{PP}\ |^2 + (1 - \frac{4(m_{l_\beta})^2}{(m_M)^2}) \left| \frac{m_M}{m_p + m_q} C_{PS} \right|^2]. \end{aligned} \tag{3.2.4}$$

We then reproduce the expression for polarization asymmetry A_{LP} (see eq. (C.12))as given by [61, 62]

$$A_{LP} = \frac{2\frac{(m_M)}{m_p+m_q} \sqrt{1 - \frac{4(m_{l_\beta})^2}{(m_M)^2}}\ \mathrm{Re}[C_{PS}(2\frac{m_{l_\beta}}{m_M} C_{AA} - \frac{(m_M)}{m_p+m_q} C_{PP})^*]}{|\ \frac{2m_{l_\beta}}{m_M} C_{AA} - \frac{(m_M)}{m_p+m_q} C_{PP}\ |^2 + (1 - \frac{4(m_{l_\beta})^2}{(m_M)^2}) \left| \frac{(m_M)}{m_p+m_q} C_{PS} \right|^2}. \tag{3.2.5}$$

Above equation can be re-arranged thus expressing A_{LP} (eq. (3.2.5)) in terms of branching fraction (eq. (3.2.4)) and C_{PS}, while assuming the coefficients to be real for practical purposes. Thus we have [61, 62],

$$A_{LP} = \frac{2\sqrt{h}}{Br[M \to l_\beta^+ l_\beta^-]} \mathrm{Re}[C_{PS} \sqrt{Br(M \to l_\beta^+ l_\beta^-) - h\,|C_{PS}|^2}], \tag{3.2.6}$$

where the constant h is defined as:

$$h = \frac{1}{8\pi}(f_M)^2 (m_{_M})^3 \tau_M (1 - \frac{4(m_{l_\beta})^2}{(m_M)^2})^{\frac{3}{2}} (\frac{(m_M)}{m_p + m_q})^2. \tag{3.2.7}$$

In SM, there is no scalar exchange particle therefore, one has $C_{PS} = C_{PP} = 0$ and $C_{AA} = \frac{Y(x_t)}{(\sin\theta_W)^2}$ arising from box and penguin diagrams, where $Y(x_t)$ is the Inami-Lim function [65] with $x_t = \frac{m_t}{M_W}$. Thus a signal of A_{LP} would be a direct indication of physics beyond SM. $R\!\!\!/_p$ framework includes scalar sparticles mediating between SM particles. Thus one can have non-zero A_{LP} in $\not{R}_p$ framework. We shall now discuss this possibility of non-zero A_{LP} in ($B_q \to l^+_\beta l^-_\beta$) in $\not{R}_p$ framework. A_{LP} arises by the presence of sneutrino and squark exchange terms. In $\not{R}_p$ MSSM the relevant effective Lagrangian is given by [60].

$$L^{eff}_{\not{R}_\mathbf{p}}\left(p\bar{q} \longrightarrow l_\beta + \bar{l}_\beta\right) = \frac{G_F}{\sqrt{2}}\left[\begin{array}{c} A^{bq}_{\beta\beta}\left(\overline{l_\beta}\gamma^\mu P_L l_\beta\right)\left(\overline{p}\gamma_\mu P_R q\right) \\ -B^{bq}_{\beta\beta}\left(\overline{l_\beta}P_R l_\beta\right)\left(\overline{p}P_L q\right) \\ -C^{bq}_{\beta\beta}\left(\overline{l_\beta}P_L l_\beta\right)\left(\overline{p}P_R q\right)\end{array}\right], \tag{3.2.8}$$

where $P_R = \frac{1+\gamma_5}{2}$; $P_L = \frac{1-\gamma_5}{2}$.The first term in eq. (3.2.8) comes from the up squark exchange and the remaining two terms come from sneutrino exchange. Here q and b are down type quarks. The dimensionless coupling constants $A^{pq}_{\beta\beta}$, $B^{pq}_{\beta\beta}$ and $C^{pq}_{\beta\beta}$ depend on the species of charged leptons and are given by [66]

$$A^{pq}_{\beta\beta} = \frac{\sqrt{2}}{G_F}\sum_{m,n,i=1}^{3} V^\dagger_{ni}V_{im}\frac{\lambda'_{\beta np}\lambda'^*_{\beta mq}}{2m^2_{\tilde{u}^c_i}}, \tag{3.2.9}$$

$$B^{pq}_{\beta\beta} = \frac{\sqrt{2}}{G_F}\sum_{i=1}^{3}\frac{2\lambda^*_{i\beta\beta}\lambda'_{iqp}}{m^2_{\tilde{\nu}_{i_L}}}, \tag{3.2.10}$$

$$C^{pq}_{\beta\beta} = \frac{\sqrt{2}}{G_F}\sum_{i=1}^{3}\frac{2\lambda_{i\beta\beta}\lambda'^*_{ipq}}{m^2_{\tilde{\nu}_{i_L}}}. \tag{3.2.11}$$

In order to calculate A_{LP}, we first re-write the relevant coefficients by comparing eq. (3.2.3) with eq. (3.2.8).

$$C_{AA} = \frac{G_F}{\sqrt{2}}\sum_{i=u,c,t} V^*_{ip}V_{iq}\left[\frac{\alpha}{2\pi}\frac{Y(x_i)}{(\sin\theta_W)^2}\right] - \frac{1}{8}\sum_{m,n,i=1}^{3} V^\dagger_{ni}V_{im}\frac{\lambda'_{\beta np}\lambda'^*_{\beta mq}}{m^2_{\tilde{u}^c_i}}, \tag{3.2.12}$$

$$C_{PS} = \frac{1}{2}\sum_{i=1}^{3}\frac{\lambda^*_{i\beta\beta}\lambda'_{iqp} - \lambda_{i\beta\beta}\lambda'^*_{ipq}}{m^2_{\tilde{\nu}_{i_L}}}, \tag{3.2.13}$$

$$C_{PP} = \frac{1}{2}\sum_{i=1}^{3}\frac{\lambda^*_{i\beta\beta}\lambda'_{iqp} + \lambda_{i\beta\beta}\lambda'^*_{ipq}}{m^2_{\tilde{\nu}_{i_L}}}. \tag{3.2.14}$$

For $M = B^0_{q'}; q = d, s$ and $\beta = \mu$; $\frac{m_{l_\beta}}{m_{B_q}} \simeq 0$, eq. (3.2.5). reduces to

$$A_{LP} = \frac{-2\,\mathrm{Re}[C_{PS}C^*_{PP}]}{|\, C_{PP}\,|^2 + |C_{PS}|^2}. \tag{3.2.15}$$

Substituting the expressions of coefficients from eqs. (3.2.9-3.2.14) into eq. (3.2.15), we get after simplifying.

$$A_{LP} = \frac{\sum\limits_{i=1}^{3}\left(\left|\frac{\lambda_{i\beta\beta}\lambda'^*_{i3q}}{m^2_{\tilde{\nu}_{i_L}}}\right|^2 - \left|\frac{\lambda^*_{i\beta\beta}\lambda'_{iq3}}{m^2_{\tilde{\nu}_{i_L}}}\right|^2\right)}{\sum\limits_{i=1}^{3}\left(\left|\frac{\lambda_{i\beta\beta}\lambda'^*_{i3q}}{m^2_{\tilde{\nu}_{i_L}}}\right|^2 + \left|\frac{\lambda^*_{i\beta\beta}\lambda'_{iq3}}{m^2_{\tilde{\nu}_{i_L}}}\right|^2\right)}. \tag{3.2.16}$$

3.3 Results and discussion

Figs. (3.3.1-3.3.4) [20], which are plotted using the data from [61, 66, 67, 68], represent our analysis of ($B^0_{s,d} \to l^+l^-; l = \mu$) and ($K^0_s \to l^+l^-; l = e, \mu$). Dominance of Yukawa coupling products ($\lambda'_{3q3}\lambda^*_{322}$ and $\lambda'_{312}\lambda^*_{311}$; $q = d(1),\ s(2)$) is assumed over ($\lambda'_{iq3}\lambda^*_{i22}$ and $\lambda'_{j12}\lambda^*_{j11},\ i = 1,\ j = 2$) in Figs. (3.3.1-3.3.4).

Fig. (3.3.1) is plotted using eq. (3.2.16). It shows that in the case of heavy mesons like $B^0_{s,d}$, A_{LP} depends only on Yukawa coupling products $\lambda'_{3q3}\lambda^*_{322}$. It shows that A_{LP} depends upon the difference between size of $\lambda'_{3q3}\lambda^*_{322}$ and $\lambda'^*_{33q}\lambda_{322}$. Since λ'_{3q3} is different in general, from λ'^*_{33q}, it is possible that $\not{R}_p$ contribution to A_{LP} may be significant from the experimental point of view.

Figs. (3.3.2 and 3.3.3) are plotted using eq. (3.2.6) and eq. (3.2.14). These figures show how A_{LP} is related to experimental bounds on branching fraction and theoretical bounds on Yukawa coupling products. Branching fraction of $B_s \to \tau^+\tau^-$ is higher than that of $B_s \to \mu^+\mu^- \sim O(100)$. So it is a favorable decay for the measurement of A_{LP} [61]. One would need to prepare about a minimum of 10^6 events of this decay ($B_s \to \tau^+\tau^-$) to search for A_{LP}. Since τ lepton is unstable, it will decay giving two neutrinos [59, 61]. Isolating these two neutrinos will be a challenge during the experiment. It may be difficult at *LHC* but Super B factories may help in future to study such decays [59].

Fig. (3.3.4) show how A_{LP} varies with the branching fraction of pure leptonic decays of K

mesons and $R\!\!\!/_p$ $MSSM$ Yukawa sneutrino couplings. The behavior is same as in the case of Figs. (3.3.2 and 3.3.3). (pure leptonic decays of B mesons). In case of leptonic decays of K mesons squark Yukawa couplings ($\lambda'_{231}\lambda'^{*}_{232}$) also contribute to A_{LP}.

Summarizing, the analysis shows that A_{LP} is mainly influenced by sneutrino-exchange terms in case of heavy mesons. The squark exchange term is suppressed by the lepton mass only in case of $B^0_{d,s}$ meson. These predictions can be tested with future experimental searches of rare B and K decays and related A_{LP} values at CERN (LHC collider) or at Super B factories as discussed in [59, 62], because of expected number of events (1 in 10^8 for $B_s \rightarrow \mu^+\mu^-$ and 1 in 10^6 for $B_s \rightarrow \tau^+\tau^-$). The comparison may indirectly verify $R\!\!\!/_p$ MSSM predictions based on dominant sneutrino exchange processes.

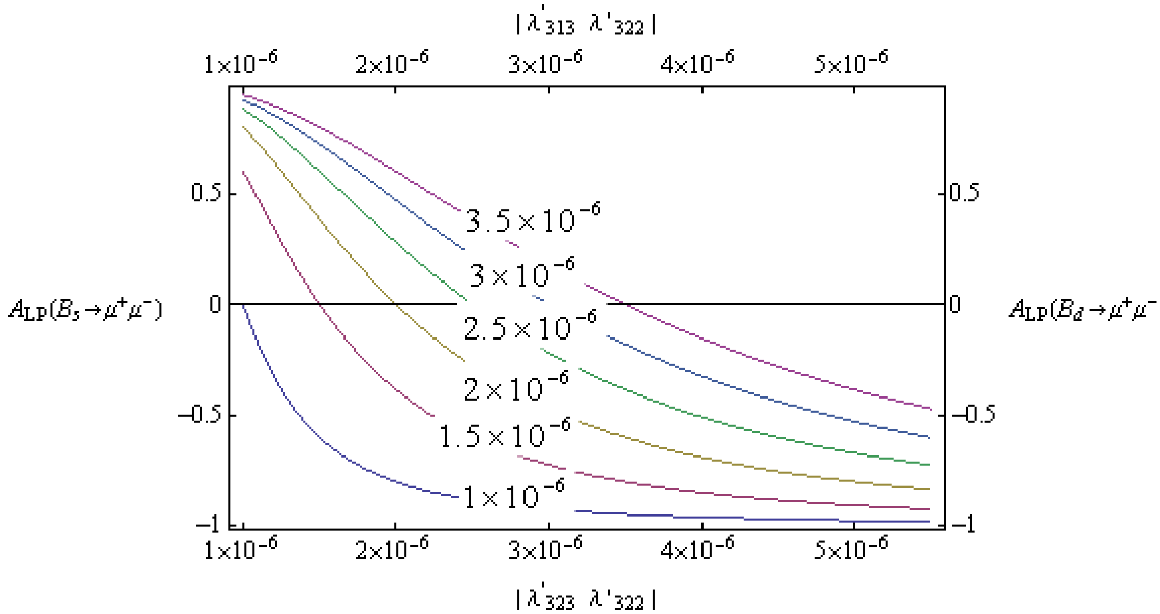


Figure 3.3.1: Dependence of A_{LP} on $\left|\lambda^*_{322}\lambda'_{323}\right|$ and $\left|\lambda^*_{322}\lambda'_{313}\right|$ at various values of $\left|\lambda_{322}\lambda'^*_{332}\right|$ and $\left|\lambda_{322}\lambda'^*_{331}\right|$ within current bounds(see section 3.2 for discussion). $\lambda'^*_{313}\lambda_{322}$ and $\lambda_{322}\lambda'^*_{332}$ are normalized to $1/(m_{\tilde{v}_{3_L}}/100GeV)^2$.

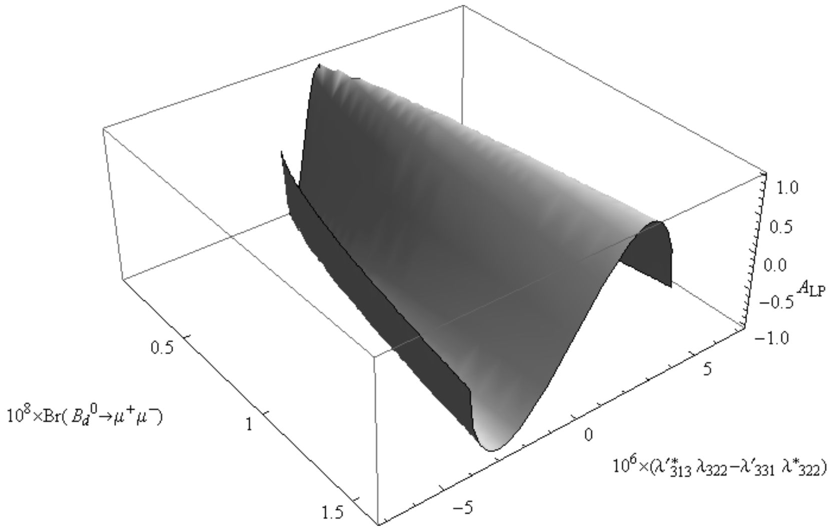


Figure 3.3.2: Dependence of A_{LP} on branching fraction of $B_d^0 \to \mu^+\mu^-$ with $\left|\lambda_{322}\lambda'^*_{313} - \lambda^*_{322}\lambda'_{331}\right|$.

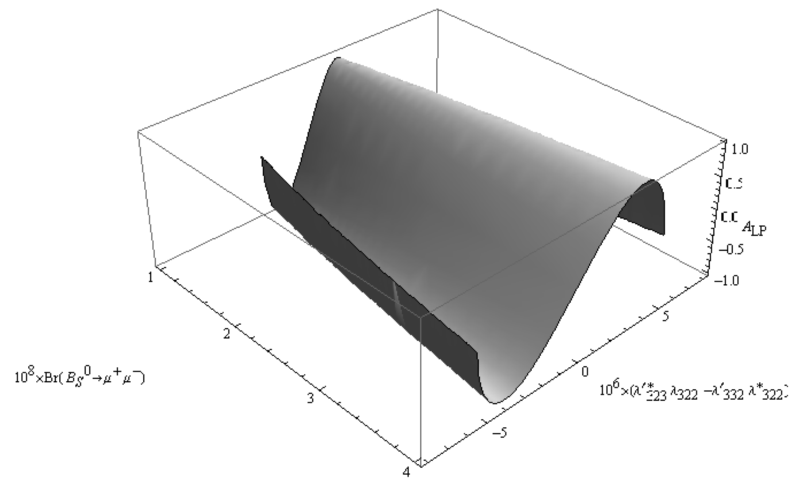
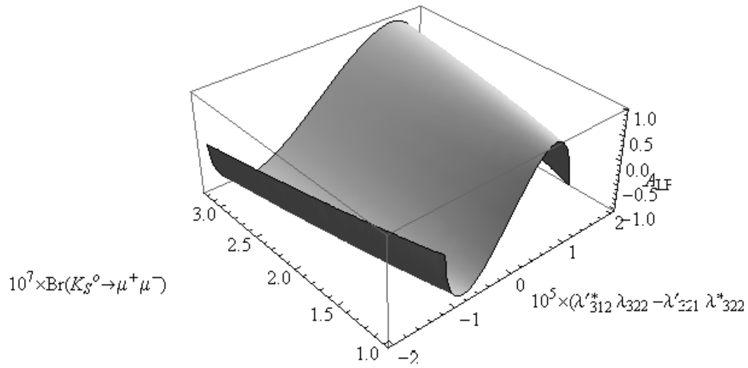


Figure 3.3.3: Dependence of A_{LP} on branching fraction of $B_s^0 \rightarrow \mu^+\mu^-$ with $\left|\lambda_{322}\lambda'^*_{323} - \lambda^*_{322}\lambda'_{332}\right|$. $\lambda'^*_{323}\lambda_{322}$ is normalized to $1/(m_{\tilde{v}_{3_L}}/100GeV)^2$.



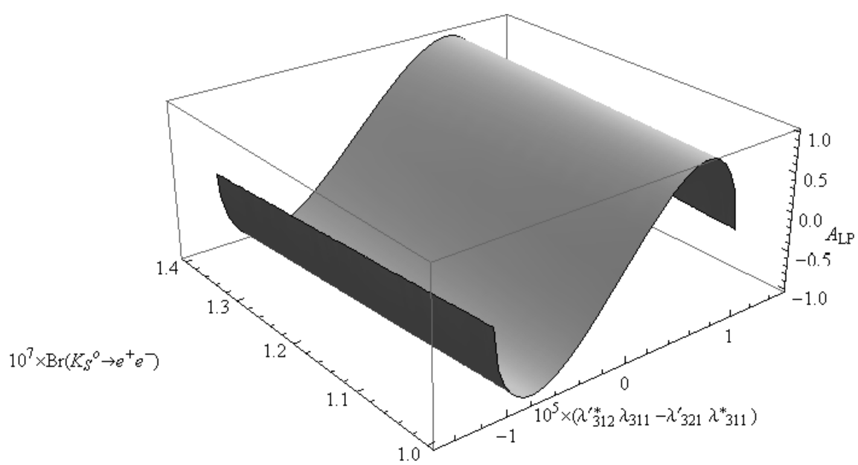


Figure 3.3.4: Dependence of A_{LP} on branching fraction of $K_S^0 \to \mu^+\mu^-$ and $K_S^0 \to e^+e^-$ on $\left|\lambda_{322}\lambda'^*_{312} - \lambda^*_{322}\lambda'_{321}\right|$ and $\left|\lambda_{311}\lambda'^*_{312} - \lambda^*_{311}\lambda'_{321}\right|$ respectively. $\lambda'^*_{312}\lambda_{322}$ and $\lambda'^*_{312}\lambda_{311}$ are normalized to $1/(m_{\tilde{\nu}_{3_L}}/100GeV)^2$ respectively.

Chapter 4

Semi-Leptonic Beauty decays $(\Delta B = 1)$ in $\not{R}_p$ MSSM

4.1 Introduction

Rare leptonic and semileptonic B decays are sensitive probe of physics beyond the SM. These decays are suppressed in SM and proceed through higher order diagrams. Such decays have also been studied in various extensions of SM [46]. Rare B decays are being searched at B factories (*Belle* and *BaBar*) [59] and also at Tevatron (*CDF* and *DO*) [30]. MSSM [60] provides a framework for these decays at tree level, where dominant SUSY particles exchange can be calculated in a simple way yielding valuable information on various phenomenological observables.

Various observables like A_{FB}, PA and CP asymmetry may also show sensitivity to NP parameters. Various studies have considered NP contribution to these observables in different models like $2HDM$ and extra dimensions and confirmed the sensitivity of these observables to the parameters of NP [54, 55].

FCNC constitutes as one important part of $\not{R}_p$SUSY phenomenology. In this chapter, we discuss the theoretical predictions of branching ratio, CP-Asymmetry, forward backward asymmetry, polarization asymmetry and double branching ratio of FCNC ($b \to sl^+l^-$) on the basis of $\not{R}_p$ in the light of current experimental data with errors [73].

4.2 Semileptonic Beauty decays

The effective Hamiltonian for the given decay process is given by [54, 74]

$$H_{eff} = \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^{*} \sum_{i=1}^{10} C_i(\mu)\ O_i(\mu) + C_Q(\mu)\ O_Q(\mu). \tag{4.2.1}$$

SM contribution in eq. (4.2.1) comes from the first set of operators ($i = 1 - 10$), while the second term describes the contribution from physics beyond SM. The scale of energy of interactions is given by μ. Wilson coefficients ($C_i(\mu); i = 1 - 10$) are given in the literature [54, 74].

In $\not{R}_p$ MSSM the relevant effective Lagrangian is given by [60]

$$L_{\not{R}_p}^{eff}\left(\bar{b} \longrightarrow \bar{s}\ l_\beta + \bar{l}_\beta\right) = \frac{G_F}{\sqrt{2}} \begin{bmatrix} A_{\beta\beta}^{bs}\left(\overline{l_\beta}\gamma^\mu P_L l_\beta\right)\left(\overline{s}\gamma_\mu P_R b\right) \\ -B_{\beta\beta}^{bs}\left(\overline{l_\beta} P_R l_\beta\right)\left(\overline{s} P_L b\right) \\ -C_{\beta\beta}^{bs}\left(\overline{l_\beta} P_L l_\beta\right)\left(\overline{s} P_R b\right) \end{bmatrix}, \beta = e, \mu, \tag{4.2.2}$$

where $P_R = \frac{1+\gamma_5}{2};\ P_L = \frac{1-\gamma_5}{2}$. The first term in eq. (4.2.2) comes from the up squark exchange and the remaining two terms come from sneutrino exchange. Here s and b denote strange and beauty down type quarks. The dimensionless coupling constants $A_{\beta\beta}^{bs}$, $B_{\beta\beta}^{bs}$ and $C_{\beta\beta}^{bs}$ depend on the species of charged leptons and are given by [66]

$$A_{\beta\beta}^{bs} = \frac{\sqrt{2}}{G_F} \sum_{m,n,i=1}^{3} \frac{V_{ni}^{\dagger} V_{im}}{2m_{\tilde{u}_i^c}^2} \lambda'_{\beta n3} \lambda'^{*}_{\beta m2}, \tag{4.2.3}$$

$$B_{\beta\beta}^{bs} = \frac{\sqrt{2}}{G_F} \sum_{i=1}^{3} \frac{2}{m_{\tilde{\nu}_{Li}}^2} \lambda^{*}_{i\beta\beta} \lambda'_{i23}, \tag{4.2.4}$$

$$C_{\beta\beta}^{bs} = \frac{\sqrt{2}}{G_F} \sum_{i=1}^{3} \frac{2}{m_{\tilde{\nu}_{Li}}^2} \lambda_{i\beta\beta} \lambda'^{*}_{i32}. \tag{4.2.5}$$

The matrix element of the decay rate ($B \to K l^+ l^-$) is given by [74]

$$M = F_S \bar{l} l + F_V (p_B)_\mu \bar{l}\gamma^\mu l + F_A (p_B)_\mu \bar{l}\gamma^\mu\gamma^5 l + F_P \bar{l}\gamma^5 l, \tag{4.2.6}$$

where $(p_B)_\mu$ is the initial momentum of B meson and $F_{S,V,A,P}$ are functions of Lorentz invariant quantities like dilepton centre of mass energy squared (s). These functions involve the Wilson coefficients ($C_i(\mu)$ and $C_{Q_i}(\mu)$) and are reproduced from [74].

We have calculated χ^2 for each set of Yukawa coupling products (sneutrino and squark exchange term) and have plotted only those sets, for which χ^2_{Br} has a minimum value in the set, where

$$\chi^2_{Br} = \sum_{i=1}^{N} \frac{(T_i - O_i)^2}{\delta(O_i)^2}.$$

N represents the number of degrees of freedom. δO_i represents errors in experimentally measured differential branching fraction O_i. T_i represents the theoretically predicted value of the branching fraction. We have converted asymmetric errors present in experimental data to symmetric errors [75]. We also divide χ^2 for each set with $N-1$, where $N = 6$ (number of experimental runs) in our case.

The Forward-Backward Asymmetry (A_{FB}) is related to asymmetric angular distribution of dilepton pair with respect to the initial meson direction of momentum in the dilepton rest frame. The A_{FB} is defined as:

$$A_{FB}(s) = \frac{\int_0^1 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta}}{\int_0^1 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta} + \int_{-1}^0 d\cos\theta \frac{d^2\Gamma}{dsd\cos\theta}}. \qquad (4.2.7)$$

In the dilepton rest frame [76, 46]

$$A_{FB}(s) = \frac{1}{128\pi^3 m_B^3 (\frac{d\Gamma}{ds})} m_l \beta(m_l, s)^2 \lambda(s) \,\mathrm{Re}(F_S F_V^*), \qquad (4.2.8)$$

where

$$\frac{d\Gamma}{ds} = \frac{1}{256\pi^3 m_B^3} \beta(m_l, s) \lambda^{\frac{1}{2}}(s) R(s) \qquad (4.2.9)$$

$$\beta(m_l, s) = (1 - \frac{4m_l^2}{s})^{\frac{1}{2}},$$

$$\lambda(s) = m_B^4 + m_K^4 + s^2 - 2sm_K^2 - 2sm_B^2 - 2m_K^2 m_B^2,$$

where

$$R(s) = |F_S|^2\, 2s\beta(m_l,s)^2 + |F_P|^2\, 2s + |F_V|^2\, \frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) + |F_A|^2\, [\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) + 8m_B^2 m_l^2]$$
$$+\mathrm{Re}(F_P F_A^*) 4m_l(m_B^2 - m_K^2 + s), \quad (4.2.10)$$

and s (invariant mass squared of the dilepton system in its rest frame) is bounded as:

$$4(m_l)^2 \leq s \leq (m_B - m_K)^2.$$

The coefficients F_V, F_A, F_S, F_P are determined by using eqs. (4.2.1, 4.2.6and A.4.1 to A.4.10)

$$F_V = \frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi}(2C_9^{eff} f^+(s) - C_7 \frac{4m_b}{m_B + m_K} f^T(s)) + \frac{1}{4} f^+(s) \sum_{i,m,n=1}^{3} V_{ni}^\dagger V_{im} \frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{m^2_{\widetilde{u^c_i}}},$$
$$F_A = \frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi}(2C_{10} f^+(s)) - \frac{1}{4} f^+(s) \sum_{i,m,n=1}^{3} V_{ni}^\dagger V_{im} \frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{m^2_{\widetilde{u^c_i}}},$$
$$F_S = -\frac{1}{2}\frac{(m_B^2 - m_K^2)}{m_b - m_s} f_o(s) \sum_{i=1}^{3} \frac{(\lambda^*_{i\beta\beta}\lambda'_{i23} + \lambda_{i\beta\beta}\lambda'^*_{i32})}{m^2_{\widetilde{\nu}_{Li}}}, \quad (4.2.11)$$
$$F_P = \frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi} 2m_l C_{10}(f^+(s) + f^-(s)) + \frac{1}{2}\frac{(m_B^2 - m_K^2)}{m_b - m_s} f_o(s) \sum_{i=1}^{3} \frac{(\lambda_{i\beta\beta}\lambda'^*_{i32} - \lambda^*_{i\beta\beta}\lambda'_{i23})}{m^2_{\widetilde{\nu}_{Li}}}.$$

Notice that in eq. (4.2.11), F_V, F_A, F_P, under H_{eff} (see eq. (4.2.1)) are split-up into two parts each carrying SM and MSSM contributions, while F_S contains only the contribution from MSSM.

We use the form factors $f^\pm(s)$, $f^T(s)$, $f_o(s)$ analytically derived in literature [74]. We use values of Wilson coefficients $C_i(m_b)$ [55, 74, 76] to leading order evaluated at the rest mass m_b of b quarks (see eqs. (D.5-D.10)). Integrated branching fraction for this process used in our calculation for comparison with the data on branching ratio is given

$$BR(B^\pm \to K^\pm l^+ l^-) = \int_{4(m_l)^2}^{(m_B - m_K)^2} \frac{d\Gamma}{ds} ds\ \tau_B, \quad (4.2.12)$$

where τ_B is the life time of B meson.

The polarization asymmetries are defined, as usual, as [55, 74]:

$$(PA)_i = \frac{\frac{d\Gamma}{ds}(\widehat{n} = -\widehat{e}_i) - \frac{d\Gamma}{ds}(\widehat{n} = \widehat{e}_i)}{\frac{d\Gamma}{ds}(\widehat{n} = -\widehat{e}_i) + \frac{d\Gamma}{ds}(\widehat{n} = \widehat{e}_i)}, \tag{4.2.13}$$

where (i=L (longitudinal), N (normal) and T (Transverse)) and $\widehat{n}$ is the spin direction of lepton l. Following three polarization unit vectors are defined in the centre of mass of the l^+l^- system and are given in literature [55, 74]:

$$\widehat{e}_L = \frac{\vec{p}_1}{|\vec{p}_1|},$$

$$\widehat{e}_N = \frac{\vec{p}_K \times \vec{p}_1}{|\vec{p}_K \times \vec{p}_1|},$$

$$\widehat{e}_T = \widehat{e}_N \times \widehat{e}_L,$$

where $\vec{p}_1$ and $\vec{p}_K$ are the three momenta of the l^- lepton and the K meson respectively. $\widehat{e}_L$ is the unit vector of lepton in the dilepton rest frame of reference. Unit vector $\widehat{e}_N$ is normal to the plane of K meson and $\widehat{e}_L$ and $\widehat{e}_T$ lie in the plane containing lepton and K meson momentum vector.

$(PA)_L$ and $(PA)_T$ have been plotted versus dilepton centre of mass energy squared and Yukawa couplings to perceive the behavior of these observables in physics beyond SM (based on R parity violating MSSM). Normal polarization asymmetry ($\widehat{e}_N$) is too small to be given any importance [55]. The expressions for polarization asymmetries are substituted from [54, 55, 74].

$$(PA)_L = \frac{\beta(m_l, s)}{R(s)}[\frac{2}{3}\lambda \,\mathrm{Re}(F_V^* F_A) - 4s \,\mathrm{Re}(F_S^* F_P) - 4m_l((m_B)^2 - (m_K)^2 - s) \,\mathrm{Re}(F_A^* F_S)], \tag{4.2.14}$$

$$(PA)_T = \frac{\pi\sqrt{\lambda(s)}}{R(s)\sqrt{s}}[m_l((m_B)^2 - (m_K)^2 - s) \,\mathrm{Re}(F_V^* F_A) + s(\beta(m_l, s))^2 \,\mathrm{Re}(F_A^* F_S) + s \,\mathrm{Re}(F_V^* F_P)], \tag{4.2.15}$$

where m_l is the mass of lepton. CP asymmetries are defined as [77]

$$A_{CP}(s) = \frac{\frac{d\Gamma(B \to K l^+ l^-)}{ds} - \frac{d\Gamma(\bar{B} \to \bar{K} l^- l^+)}{ds}}{\frac{d\Gamma(B \to K l^+ l^-)}{ds} + \frac{d\Gamma(\bar{B} \to \bar{K} l^- l^+)}{ds}}. \tag{4.2.16}$$

We have also computed average CP-asymmetries, which are defined as

$$< A_{CP} >= \frac{\int_{4(m_l)^2}^{(m_B-m_K)^2} A_{CP}(s)\frac{dBr}{ds}ds}{\int_{4(m_l)^2}^{(m_B-m_K)^2} \frac{dBr}{ds}ds}. \tag{4.2.17}$$

In SM, C_9^{eff} becomes complex due to non-negligible terms induced in A and B [78]. Its complex nature is responsible for CP-asymmetry. C_7 and C_{10} being real do not contribute to CP-asymmetry. $R\!\!\!/_p$ Yukawa coupling products can be imaginary, so we have

$$\begin{aligned}\frac{d\Gamma(\bar{B}\to\bar{K}l^-l^+)}{ds} &= \frac{1}{256\pi^3 m_B^3}\beta(m_l,s)\lambda^{\frac{1}{2}}(s)\{\left|\bar{F}_S\right|^2 2s\beta(m_l,s)^2+\left|\bar{F}_P\right|^2 2s \\ &+\left|\bar{F}_V\right|^2\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s})+\left|\bar{F}_A\right|^2[\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s})+8m_B^2m_l^2] \\ &+\mathrm{Re}(\bar{F}_P^*\bar{F}_A)4m_l(m_B^2-m_K^2+s)\},\end{aligned} \tag{4.2.18}$$

$\bar{F}_V,\ \bar{F}_A,\ \bar{F}_P,\ \bar{F}_S$ are given by eq. (E.6)) by taking complex conjugate at appropriate places

$$A_{CP}=\lambda(s)(1+\frac{2m_l^2}{s})(|F_V|^2-\left|\bar{F}_V\right|^2)(D(s))^{-1}, \tag{4.2.19}$$

where

$$\begin{aligned}D(s) &= 6(|F_S|^2 2s\beta(m_l,s)^2+|F_P|^2 2s+|F_A|^2[\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s})+8m_B^2m_l]+ \\ &\mathrm{Re}(F_PF_A^*)4m_l(m_B^2-m_K^2+s)+\frac{\lambda(s)}{6}(1+\frac{2m_l^2}{s})(|F_V|^2+\left|\bar{F}_V\right|^2)).\end{aligned} \tag{4.2.20}$$

We have calculated average CP-asymmetries to determine its dependence on $R\!\!\!/_p$ Yukawa coupling products.

We also study the double branching ratio [79]

$$R_H=\frac{BR(B^\pm\to K^\pm\mu^+\mu^-)}{BR(B^\pm\to K^\pm e^+e^-)}=\frac{\int_{4(m_\mu)^2}^{(m_B-m_K)^2}\frac{d\Gamma}{ds}ds}{\int_{4(m_e)^2}^{(m_B-m_K)^2}\frac{d\Gamma}{ds}ds}. \tag{4.2.21}$$

Eqs. (4.2.8, 4.2.9, 4.2.10, 4.2.11 and 4.2.21) show that R_H is a function of R-parity Yukawa coupling products and can be used to study bounds given the values of these ratios from experimentally measured parameters. In the next section, some results on asymmetries in this process which are computed from the eqs. (4.2.14-4.2.19) and will discuss their comparisons with present and future experiments.

4.3 Results and discussion

Our analysis of the FCNC process $B \to K\mu^+\mu^-$ within $\not{R}_p$ $MSSM$ model has been summed up in Figs. (4.3.1-4.3.8) [20]. These figures are plotted using experimental data from [73, 67]. They compare experimental results on branching fraction and A_{FB} with the theoretical predictions made within $\not{R}_p$ $MSSM$. These figures correspond to the set of Yukawa coupling products (sneutrino and squark exchange terms) for which χ^2_{Br} has a minimum value.

Fig. (4.3.2), compares differential branching fraction, calculated within both $\not{R}_p$ $MSSM$ and SM, with the experimental data. In the case of SM (shown by the line 1), one obtains a somewhat poor value of χ^2 $((\chi^2_{Br})_{SM}/5 = 2.1)$, which shows the lack of sufficient experimental data for comparison. As discussed in section 1.4 $\not{R}_p$ MSSM has parameters $(\lambda^*_{322}\lambda'_{323}, \lambda_{322}\lambda'^*_{332}, \lambda'_{233}\lambda'^*_{232})$ given by eq. (1.4.1). For the particular ordered set of Yukawa coupling products $(\lambda^*_{322}\lambda'_{323}, \lambda_{322}\lambda'^*_{332}, \lambda'_{233}\lambda'^*_{232})$, one obtains ($\chi^2_{Br}/5 = 0.76$ and 0.75). This shows that $\not{R}_p$ $MSSM$ fits the experimental data quite nicely for this particular set. The total branching fraction for SM is calculated to be 5.22×10^{-7}. In case of $\not{R}_p$ $MSSM$ total branching fraction values calculated corresponding to two figures are (4.19×10^{-7} and 4.22×10^{-7}) respectively.

Theoretical prediction of A_{FB} in SM and $\not{R}_p$ $MSSM$ is compared with the experimental data in Figs. (4.3.3). In SM case, one obtains $\chi^2_{SM}/5 = 1$ showing that the experimental data for A_{FB} fits nicely with SM prediction. Also in $\not{R}_p$ $MSSM$ case, one obtains $\chi^2_{Br}/5 = 1.2$ and 0.96 indicating a good fit of the experimental data. One may conclude that $\not{R}_p$ $MSSM$ predicts a vanishingly small A_{FB} in the case of $B \to Kl\bar{l}$, which is in good agreement with the given experimental data and prediction of SM.

Fig. (4.3.4) studies the variation of differential branching fraction w.r.t. centre of mass energy squared at several values of masses of squarks ($800GeV \leq m_{\widetilde{u}} \leq 2TeV$) and sneutrinos ($400GeV \leq m_{\widetilde{\nu}} \leq 1TeV$) and fixed values of $\lambda^*_{322}\lambda'_{323} = -3 \times 10^{-5}$ and $\lambda'_{233}\lambda'^*_{232} = 3 \times 10^{-6}$. The curves with sneutrinos with mass 1 TeV and squarks with mass 1.2 TeV fit nicely with given experimental data as they have lowest χ^2_{Br} within a set of values i.e., $\chi^2_{Br}/5 = 0.8$ for $m_{\tilde{\nu}} = 1TeV$ and $\chi^2_{Br}/5 = 0.85$ for $m_{\tilde{u}} = 1.2TeV$. This hints at the possible range of values for the masses of sparticles, which may be explored at LHC [81].

$(PA)_L$ and $(PA)_T$ have also been studied in Figs. (4.3.5 and 4.3.6) respectively, showing that in $\not{R}_p$ $MSSM$ case, we do not get any new results different from SM. Fig. (4.3.7) presents a comparison of theoretical prediction of A_{CP}, within SM and $\not{R}_p$ $MSSM$, with the experimental data. The experimentally measured value of A_{CP} is $(0.04 \pm 0.1 \pm 0.02)$ [73]. This Fig. (4.3.7) yields average values of A_{CP} as -1% and -2.2% for different sets of Yukawa coupling products. This theoretical prediction is compatible with the experimentally measured value within 1σ error(see Fig. (4.3.7). The presence of a wide error in experimental measurement indictes the requirement of more data.

In Fig. (4.3.8), we plot the region which shows the allowed values of Yukawa coupling products constrained by the observed values of double branching ratio $(1.03 \pm 0.19 \pm 0.06)$ [73] within 1σ and 2σ errors. Single coupling dominance has been assumed in these graphs. We have also assumed real Yukawa couplings for simplicity. We have used the CP-violating phase $\delta(\lambda_t) = 45^o$ and 60^o in Figs. (4.3.2-4.3.7), which is within 1σ of the mean value (77^o).where

$$\lambda_t = \frac{V_{ub}^* V_{us}}{V_{tb}^* V_{ts}}$$

Squark Yukawa coupling products have been assumed to be complex and sneutrino exchange Yukawa coupling products are real as their phases have no significant contribution to the differential branching fraction and forward backward asymmetry. The complex part of sneutrino Yukawa coupling products is suppressed by lepton mass term (coefficient of $F_P F_A^*$ in eq. (4.2.10)).

Summarizing, we have studied the decay processes $(B^{\pm} \to K^{\pm} l^+ l^-)$ in both SM and $\not{R}_p$ MSSM. Theoretical predictions for various observables like A_{FB}, PAs and A_{CP} have been compared with the experimental data and it is concluded that the predictions made by $\not{R}_p$ MSSM agree well with the given experimental data and SM [73]. These observables may be measured at future experiments at Super B factory [51] and may provide a good probe for physics beyond SM.

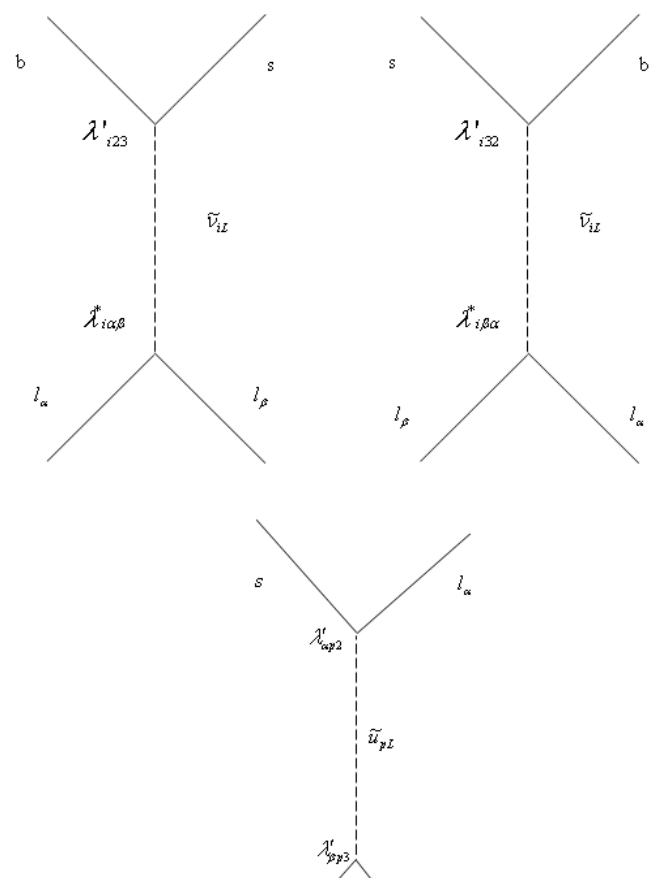


Figure 4.3.1: Tree-Level diagrams contributing to charged B decays, $B^\pm \to K^\pm l^+ l^- (\bar{b} \to \bar{s} l^+ l^-; l = e, \mu)$.

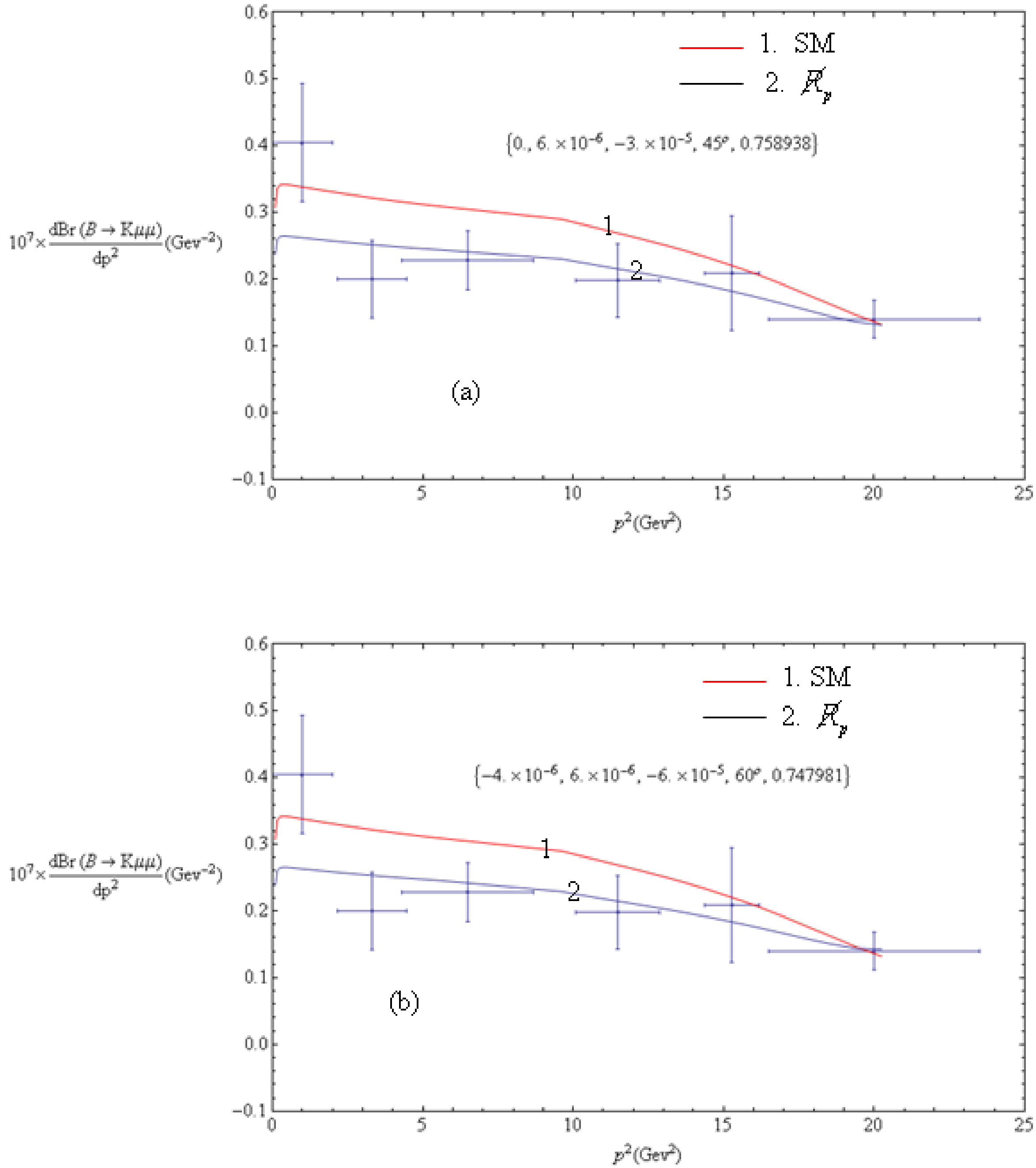
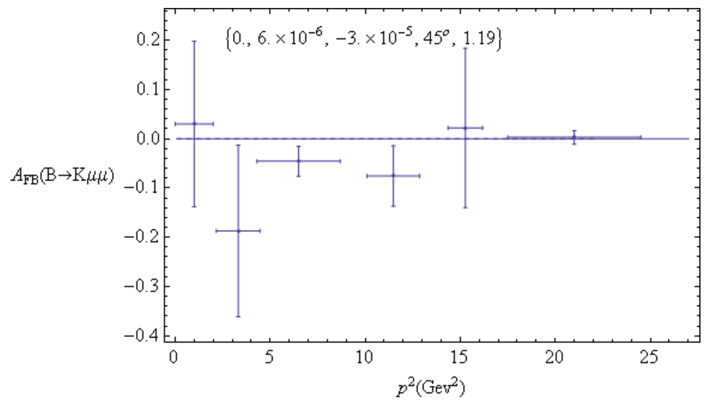


Figure 4.3.2: Differential branching fraction for $B^{\pm} \to K^{\pm} l^{+} l^{-}$as a function of p^2. The curves show the theoretical calculation (2) from R-parity violating MSSM and also from SM(1). Error bars denote experimental data. The ordered set $\{\lambda^{*}_{322}\lambda'_{323}, \lambda_{322}\lambda'^{*}_{332}, \lambda'_{233}\lambda'^{*}_{232}, \delta(\lambda_t),$ $\chi^2_{Br}/5\}$ represents the case for minimum χ^2_{Br}.$\lambda^{*}_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{v}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^{*}_{232}$ is expressed in the units of $1/(m\ _{\tilde{u}_{i_L}}/100GeV)^2$.



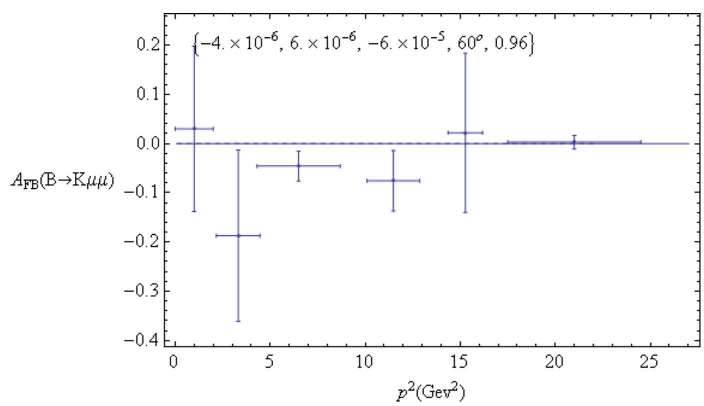


Figure 4.3.3: Forward-backward asymmetry for $B^{\pm} \rightarrow K^{\pm}l^{+}l^{-}$as a function of p^2. The curves show the theoretical calculation (solid) from R-parity violating MSSM and also from SM (dotted). These two curves overlap showing agreement of two models. Error bars denote experimental data. The ordered set $\{\lambda^*_{322}\lambda'_{323}, \lambda_{322}\lambda'^*_{332}, \lambda'_{233}\lambda'^*_{232}, \delta(\lambda_t), \chi^2_{Br}/5\}$ represents the case for minimum χ^2_{Br}. $\lambda^*_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{\nu}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^*_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.

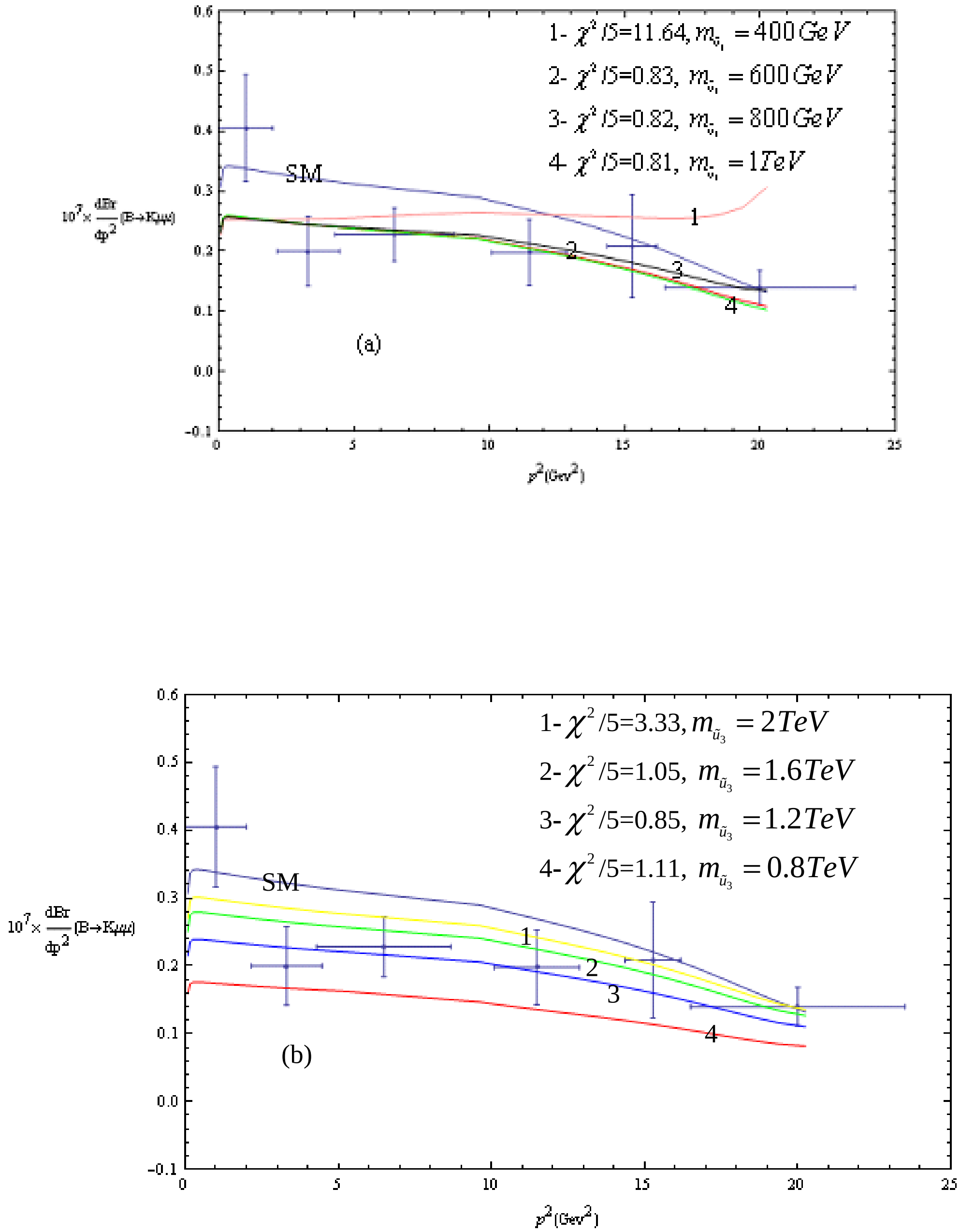
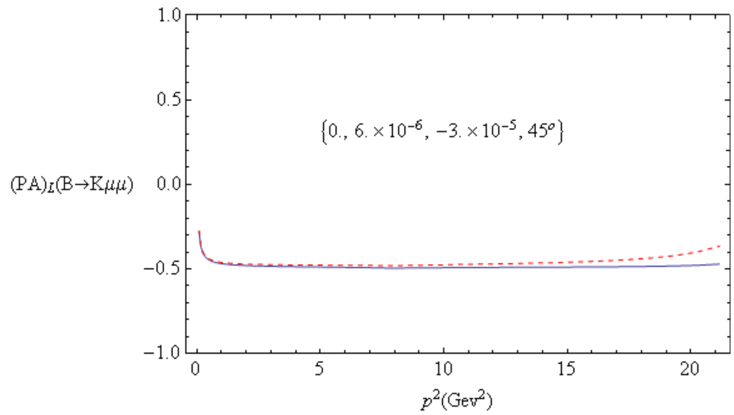


Figure 4.3.4: The curves show the variation of $\frac{dBr}{ds}$ $(B^{\pm} \to K^{\pm}\mu^{+}\mu^{-})$ w.r.t. p^2 at several values of sparticle masses and fixed values of $\lambda^{*}_{322}\lambda'_{323}$ and $\lambda'_{233}\lambda'^{*}_{232}$. Running sneutrino mass at $\lambda^{*}_{322}\lambda'_{323} = -3 \times 10^{-5}$. Running squark mass $\lambda'_{233}\lambda'^{*}_{232} = 3 \times 10^{-6}$. Error bars denote experimental data.



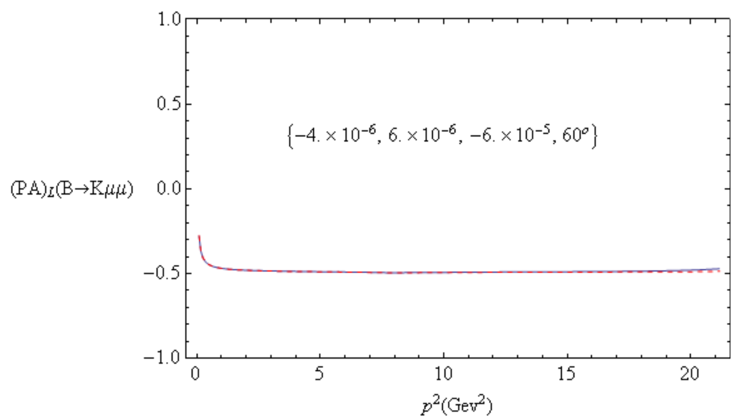
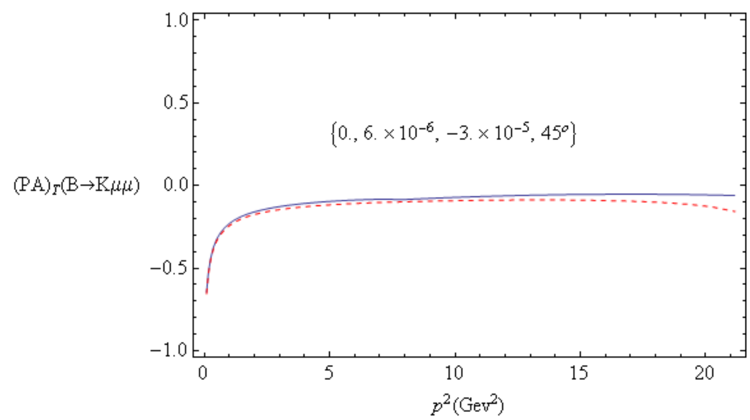


Figure 4.3.5: Longitudinal polalrization asymmetry for $B^{\pm} \to K^{\pm}l^{+}l^{-}$as a function of p^2. The curves show the theoretical calculation (solid) from R-parity violating MSSM and also from SM (dotted). The ordered set $\{\lambda^*_{322}\lambda'_{323}, \lambda_{322}\lambda'^*_{332}, \lambda'_{233}\lambda'^*_{232}, \delta(\lambda_t),\ \chi^2_{Br}/5\}$ represents the case for minimum χ^2_{Br}. $\lambda^*_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{\nu}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^*_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.



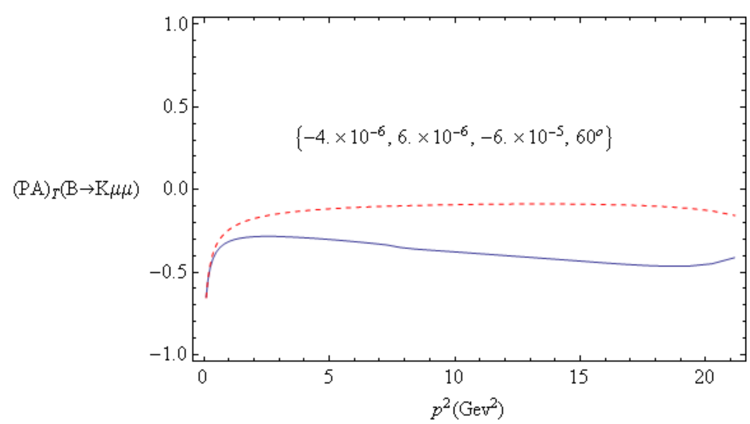


Figure 4.3.6: Transverse polarization asymmetry for $B^{\pm} \to K^{\pm} l^{+} l^{-}$as a function of p^2. The curves show the theoretical calculation (solid) from R-parity violating MSSM and also from SM (dotted). The ordered set $\{\lambda^{*}_{322}\lambda'_{323}, \lambda_{322}\lambda'^{*}_{332}, \lambda'_{233}\lambda'^{*}_{232}, \delta(\lambda_t),\ \chi^2_{Br}/5\}$ represents the case for minimum χ^2_{Br}. $\lambda^{*}_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{v}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^{*}_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.

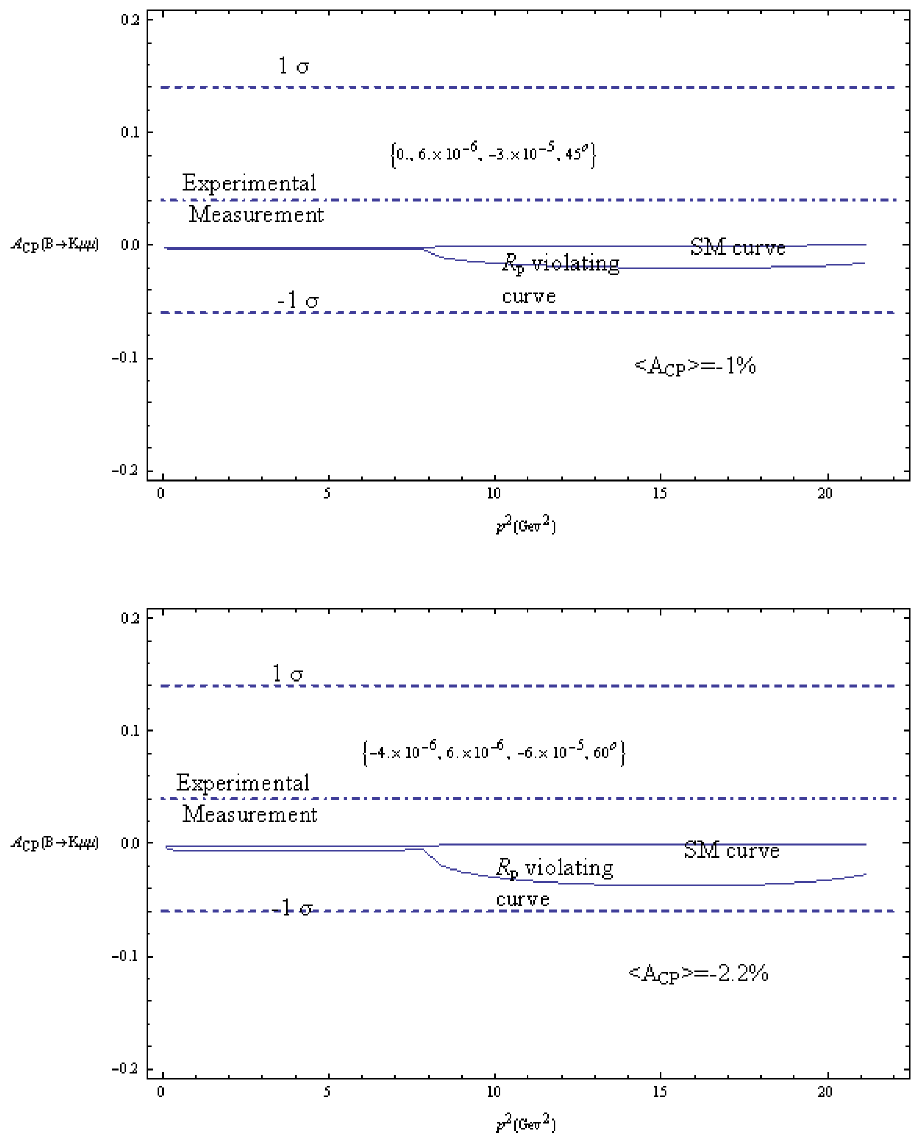


Figure 4.3.7: CP-asymmetry for $B^{\pm} \rightarrow K^{\pm}l^{+}l^{-}$as a function of p^2. The curves show the theoretical calculation (solid) from R-parity violating MSSM and also from SM (dotted). The ordered set $\{\lambda^{*}_{322}\lambda'_{323}, \lambda_{322}\lambda'^{*}_{332}, \lambda'_{233}\lambda'^{*}_{232}, \delta(\lambda_t),\ \chi^2_{Br}/5\}$ represents the case for minimum χ^2_{Br}. $\lambda^{*}_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{\nu}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^{*}_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.

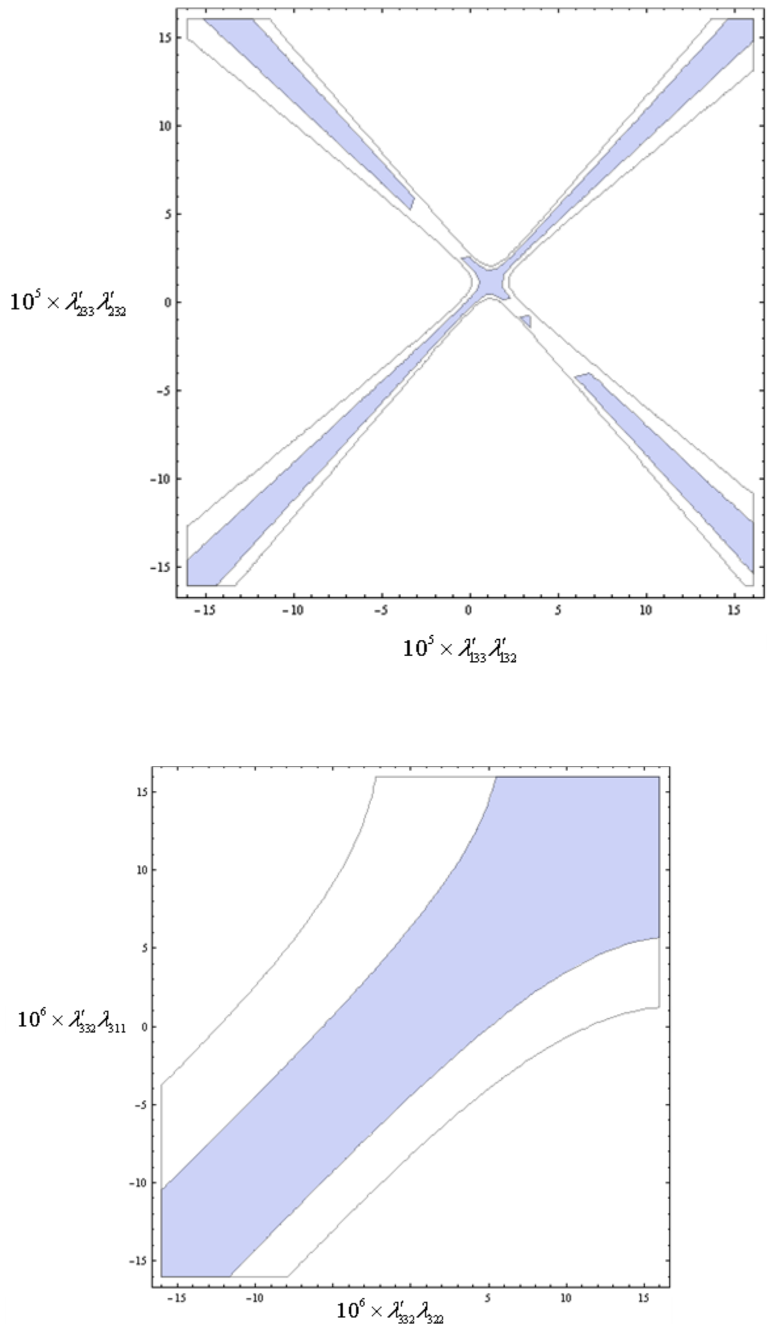


Figure 4.3.8: The shaded and un shaded outlined region represents the allowed values of Yukawa coupling products for which double branching ratio is within 1σ and 2σ experimental errors. Single coupling dominance is assumed. Yukawa couplings have been assumed real for simplicity in this analysis. $\lambda^*_{322}\lambda'_{323}$ is expressed in the units of $1/(m_{\tilde{v}_{i_L}}/100GeV)^2$. $\lambda'_{233}\lambda'^*_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$.

Chapter 5

Summary and conclusion

This thesis present the studies of leptonic rare meson decays within the framework of $\not{R}_p MSSM$ phenomenology. Various physical observables like forward backward asymmetry and polarization asymmetry in such decay processes have also been studied.

In **chapter 1**, SM is introduced with some of its inadequacies like hierarchy problem and fermion mass problem. Supersymmetry presents a solution for hierarchy problem. In addition, it predicts the existence of new particles as partners of SM particles. These particles must have the same mass as of SM particles if SUSY is not broken. As this is not the case indicating that it must be broken, if present in nature. $MSSM$ is discussed along with the idea of superpotential. A discussion on R-parity conservation is also made in this chapter. Phenomenological consequences of R-parity violation are also explored. A brief discussion of searches for signals of R-parity violation is made at the end of chapter.

Chapter 2 focuses on the study of rare semileptonic kaon decays ($K^{\pm} \to \pi^{\pm} l_{\alpha}^{+} l_{\beta}^{-}; \alpha, \beta = e, \mu$) within $\not{R}_p$ $MSSM$. Bounds on Yukawa couplings ($\lambda_{i\alpha\beta}\lambda'^{*}_{i21}, \lambda'^{*}_{\alpha i1}\lambda'_{\beta i2}$) have been calculated (see Table. (2.3.1)) under the assumption that only one product combination is non zero. Tables. (2.3.1and 2.3.2) have been obtained by using inclusive decay mode i.e., using eq. (2.2.14) [20]. The Yukawa couplings $\lambda_{i\alpha\beta}\lambda'^{*}_{i21}$ have been normalized to $1/(m_{\widetilde{\nu}_{i_L}}/100GeV)^2$ and $\lambda'_{\alpha i1}\lambda'^{*}_{\beta i2}$ to $1/(m_{\widetilde{u}_{i_L}}/100GeV)^2$. Table. (2.3.1) shows that the bounds on squark exchange products $\lambda'_{\alpha i1}\lambda'^{*}_{\beta i2}$ are larger $\sim O(10)$ only, while the bounds on sneutrino Yukawa coupling products $\lambda_{i\alpha\beta}\lambda'^{*}_{i21}$ are larger$\sim O(10^2 - 10^3)$.

The order of magnitude relation between form factors $F_S,\ F_V,\ F_P,\ F_A\ \sim\ O(10^{-9})$ for the maximum A_{FB} in these decay processes ($K^{\pm} \rightarrow \pi^{\pm} l_{\alpha}^{+} l_{\beta}^{-}; \alpha, \beta = e, \mu$). The maximum A_{FB} calculated is of $O(10^{-3})$ for $K^{+} \longrightarrow \pi^{+} e^{+} e^{-}$, which is vanishingly small and hence not measurable. While, the maximum A_{FB} is $O(10^{-1})$ for $K^{+} \longrightarrow \pi^{+} \mu^{+} \mu^{-}$ [20], which may be possible to be measured at future experiment looking for $K^{+} \longrightarrow \pi^{+} \mu^{+} \mu^{-}$.

In **chapter 3**, rare leptonic decays of K and B mesons have been considered and longitudinal polarization asymmetry, which is zero in SM have also been calculated for these decay processes. It is found to be suppressed in the case of leptonic decay of B meson due to its large mass as compared to that of leptons.

Dominance of Yukawa coupling products ($\lambda'_{3q'3}\lambda^{*}_{322}, \lambda'_{312}\lambda^{*}_{322}$), expressed in the units of $(100GeV/m_{\widetilde{\nu}_{3_L}})^2$ has been assumed over $i = 1, 2$ values in Figs. (3.3.1-3.3.4). Fig. (3.3.1) shows dependence of A_{LP} on given limits of Yukawa coupling products. It also shows that A_{LP} is non-zero for $B \rightarrow \mu^{+}\mu^{-}$ only if size of $\lambda'_{3q'3}$is different from $\lambda'^{*}_{33q'}$. Figs. (3.3.2 and 3.3.3) study the dependence of A_{LP} on branching fraction and Yukawa coupling products. $B_{s,d} \rightarrow \tau^{+}\tau^{-}$ has overall high branching fraction as predicted by SM i.e., 8×10^{-7} [61]. So it is an ideal decay process for the measurement of A_{LP} at Super B factories [59].

Chapter 4 deals with rare semi-leptonic meson decays ($B^{\pm} \rightarrow K^{\pm}\mu^{+}\mu^{-}$). It actually compares theoretically calculated results (SM and $\not{R}_p$ $MSSM$) with given experimental results [73].

Fig.(4.3.2) shows a comparison between $\not{R}_p$ $MSSM$ and SM differential branching fraction of ($B^{\pm} \rightarrow K^{\pm}\mu^{+}\mu^{-}$). This figure shows that $\not{R}_p$ $MSSM$ fits the experimental data nicely, while the SM does not give a very good fit. Fig. (4.3.3) gives a comparison of experimental data on A_{FB} with its theoretical prediction in SM and $\not{R}_p$ $MSSM$. Both SM and $\not{R}_p$ $MSSM$ give a good fit of experimental data, which predicts vanishingly small A_{FB} in ($B^{\pm} \rightarrow K^{\pm}\mu^{+}\mu^{-}$). Fig. (4.3.4) compares the same experimental data, as in Fig. (4.3.2), at various values of masses of squarks and sneutrinos. The curves with sneutrinos with mass 1 TeV and squarks with mass 1.2 TeV fit the experimental data nicely and are the candidates to look for Supersymmetry signals in ($B^{\pm} \rightarrow K^{\pm}\mu^{+}\mu^{-}$). Fig. (4.3.7) shows an analysis of A_{CP} in both SM and $\not{R}_p$ $MSSM$. This Fig. (4.3.7) shows that the theoretical prediction of $\not{R}_p$ $MSSM$ is much compatible to the experimentally measured value within $\pm 1\sigma$ error.

The decay processes ($B^{\pm} \rightarrow K^{\pm}l^{+}l^{-}$ and $B_s^0 \rightarrow l^{+}l^{-};\ l = e,\ \mu$) depend on the same set of

R-parity violating Yukawa couplings, because they represent the same flavor changing neutral current on the quark level i.e., ($b \to sl^+l^-$). Branching fraction of these two processes are compared in Figs. (5.1.1 and 5.1.2). This comparison is further illustrated in Table (5.1.1). Table (5.1.1) shows that SM contribution to branching fraction of $B_s^0 \to e^+e^-$ is very small as compared to current experimental limits [84] and is comparable to contribution from Squark Yukawa couplings ($\lambda'_{233}\lambda'^{*}_{232}$). While contribution from sneutrino Yukawa couplings ($\lambda^*_{311}\lambda'_{323}, \lambda_{311}\lambda'^{*}_{332}$) is one order less only than the current experimental limits. SM contribution to branching fraction of $B_s^0 \to \mu^+\mu^-$ is one order smaller only than the experimental limits [84]. While contribution from sneutrino Yukawa couplings ($\lambda^*_{311}\lambda'_{323}, \lambda_{311}\lambda'^{*}_{332}$) is comparable to the current experimental limits.

SM contribution to semileptonic decay process ($B^\pm \to K^\pm l^+l^-; l = e, \mu$) is within the experimental limit i.e. $\pm 1\sigma$, which is also true for R-parity Yukawa couplings ($\lambda'_{233}\lambda'^{*}_{232},\ \lambda^*_{322}\lambda'_{323},\ \lambda_{322}\lambda'^{*}_{332}$).

These observations signify the importance of role of both type of Yukawa coupling products in the study of these different modes of decay. These predictions regarding ($B^\pm \to K^\pm l^+l^-$ and $B^\pm \to l^+l^-(l = e,\ \mu)$) may be checked in the currently running **LHC**, B **factories** and future projects like **Super B factories**.

5.1 Future possible work

In this thesis, we have discussed the prospects for new physics in Flavor changing neutral current (FCNC) involving some rare leptonic decays of mesons (K and B). Our discussion is limited to $\not{R}_p$ MSSM framework only. One can extend the analysis discussed in previous chapters for the following new processes:

1. $D_s^- \to K^*(892)^-\mu^+\mu^-$.

Above mentioned process is described by $c \to u\mu^+\mu^-$. The experimental bound on this decay process is given as [67]

$$Br(D_s^- \to K^*(892)^-\mu^+\mu^-) < 1.4 \times 10^{-3}$$

2. $B^- \to K^*(892)^-l^+l^-$.

On quark level, this process is described by $b \to sl^+l^-$. The experimentally measured branching fraction on this decay process is given as [67]

$$Br(B \to K^*(892)e^+e^-) = (8 \pm 8) \times 10^{-7}$$
$$Br(B \to K^*(892)\mu^+\mu^-) = (8^{+6}_{-4}) \times 10^{-7}$$

3. $B^0 \to K^*(892)^0 l^+ l^-$.

This process can be described by $b \to sl^+l^-$. The experimentally measured branching fraction on this decay process is given as [67]

$$Br(B^0 \to K^*(892)^0 e^+e^-) = (1.04^{+0.35}_{-0.31}) \times 10^{-6}$$
$$Br(B^0 \to K^*(892)^0 \mu^+\mu^-) = (1.10^{+0.29}_{-0.26}) \times 10^{-6}$$

These decay processes can be studied within $\not{R}_p$ $MSSM$. They proceed through tree diagrams, mediated by sparticles $(\widetilde{\nu}, \widetilde{u}, \widetilde{d})$.

Observables related to these decay processes like A_{FB} and polarization asymmetries can also be studied. An interesting study would be to compare the contribution of loop diagrams of SM and MSSM R_p conservation (box, penguin and radiative) with $\not{R}_p$ $MSSM$ tree diagrams. It would be also interesting to study how the parameters from different model, for e.g. $(m_{H^\pm}, \tan\beta)$ from R_p conservation [80] and $(\lambda_{i\alpha\beta}\lambda'^*_{i21}, \lambda'^*_{\alpha i1}\lambda'_{\beta i2})$ from $\not{R}_p$ $MSSM$ are related to each other given the observed limits on the branching fraction of said decay processes.

Another interesting issue is to study supersymmetry breaking [82] with regards to these decay processes. For example, the parameters related to gauge mediated symmetry breaking are M, N, Λ, $\tan\beta$ and sign(μ), where N is the number of messenger superfields, M is the messenger scale, and Λ is the is SUSY breaking scale. Again, these parameters can be studied in relation to $\not{R}_p$ $MSSM$ Yukawa couplings.

Decay Process	SM contribution to BF	Expt on BF[85]	BF in the presence of $\lambda'_{233}\lambda'^{*}_{\alpha 3\alpha}$ ($\alpha = 1,2$)	BF in the presence of $\lambda^{*}_{3\alpha\alpha}\lambda'_{323}$ ($\alpha = 1,2$)
	(Branching Fraction)		$1/(m_{\tilde{u}_{3_L}}/100GeV)^2$	$1/(m_{\tilde{\nu}_{3_L}}/100GeV)^2$
$B_S^0 \to e^+e^-$	9.1×10^{-14}	$<2.8 \times 10^{-7}$	$< 7 \times 10^{-14}$	$<5.8 \times 10^{-6}$
$B_S^0 \to \mu^+\mu^-$	3.78×10^{-9}	$<4.7 \times 10^{-8}$	$< 3.2 \times 10^{-9}$	$<4.7 \times 10^{-8}$
$B \to K e^+e^-$	5.22×10^{-7}	$(5.5 \pm 0.7) \times 10^{-7}$	$\leq 4.9 \times 10^{-7}$	$\leq 2.83 \times 10^{-7}$
$B \to K \mu^+\mu^-$	5.22×10^{-7}	$(5.2 \pm 0.7) \times 10^{-7}$	$\leq 3.9 \times 10^{-7}$	$\leq 2.75 \times 10^{-7}$

Table 5.1.1: Comparison of contribution to branching fraction of B meson from SM and R-parity violating SUSY Yukawa couplings.

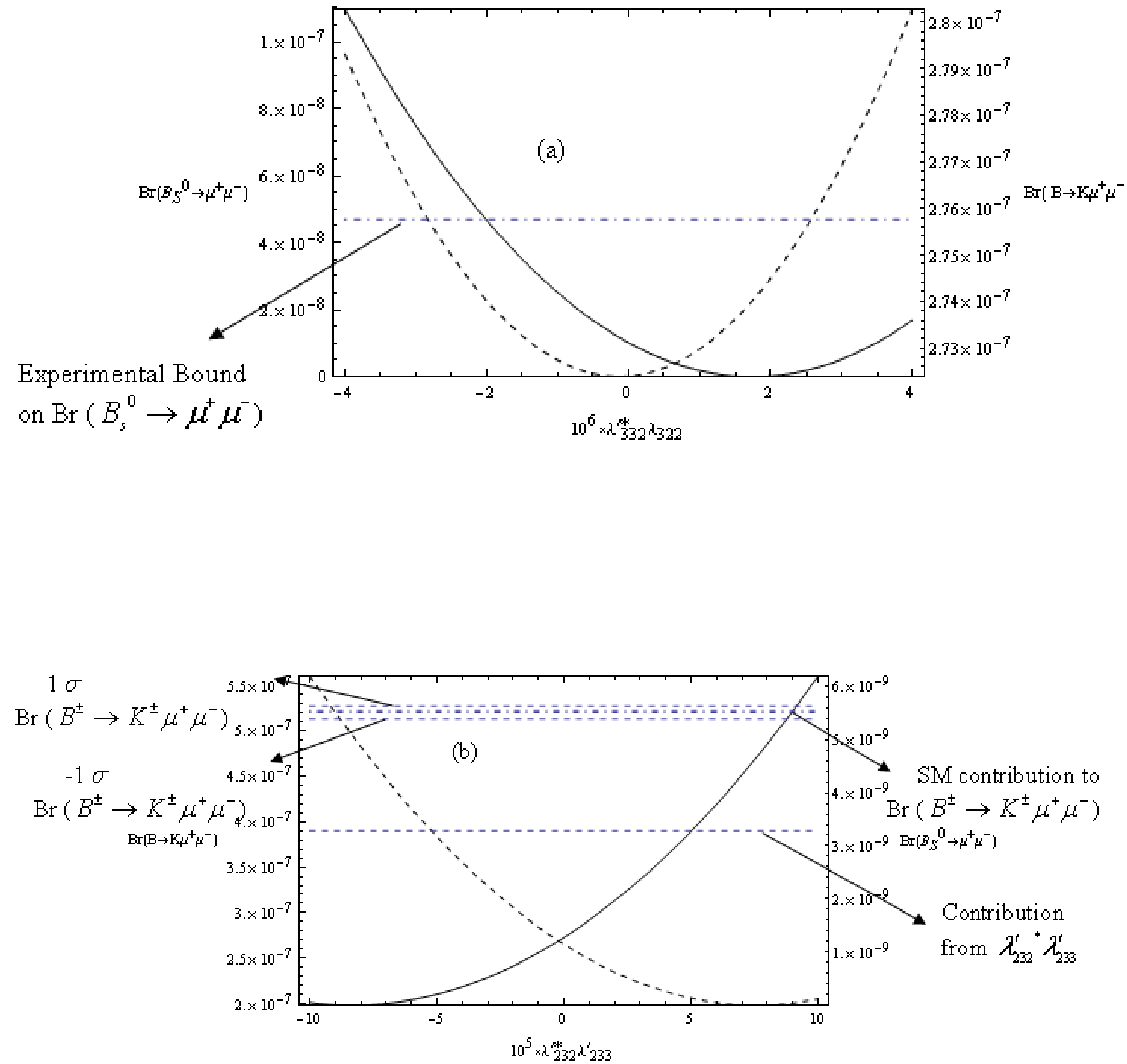


Figure 5.1.1: Comparison of branching fraction of decay process $B^\pm \to K^\pm\mu^+\mu^-$ (solid curve) with $B_s^0 \to \mu^+\mu^-$ (dashed curve) (a) as a function of sneutrino Yukawa couplings only as a function of (b) squark Yukawa couplings only. Dash-dot line (b) represents Experimental bound on $B^\pm \to K^\pm\mu^+\mu^-$. $\lambda'_{233}\lambda'^*_{232}$ is expressed in the units of $1/(m_{\tilde{u}_{i_L}}/100GeV)^2$. $\lambda'^*_{332}\lambda_{322}$ is normalized to $1/(m_{\tilde{\nu}_{3_L}}/100GeV)^2$.

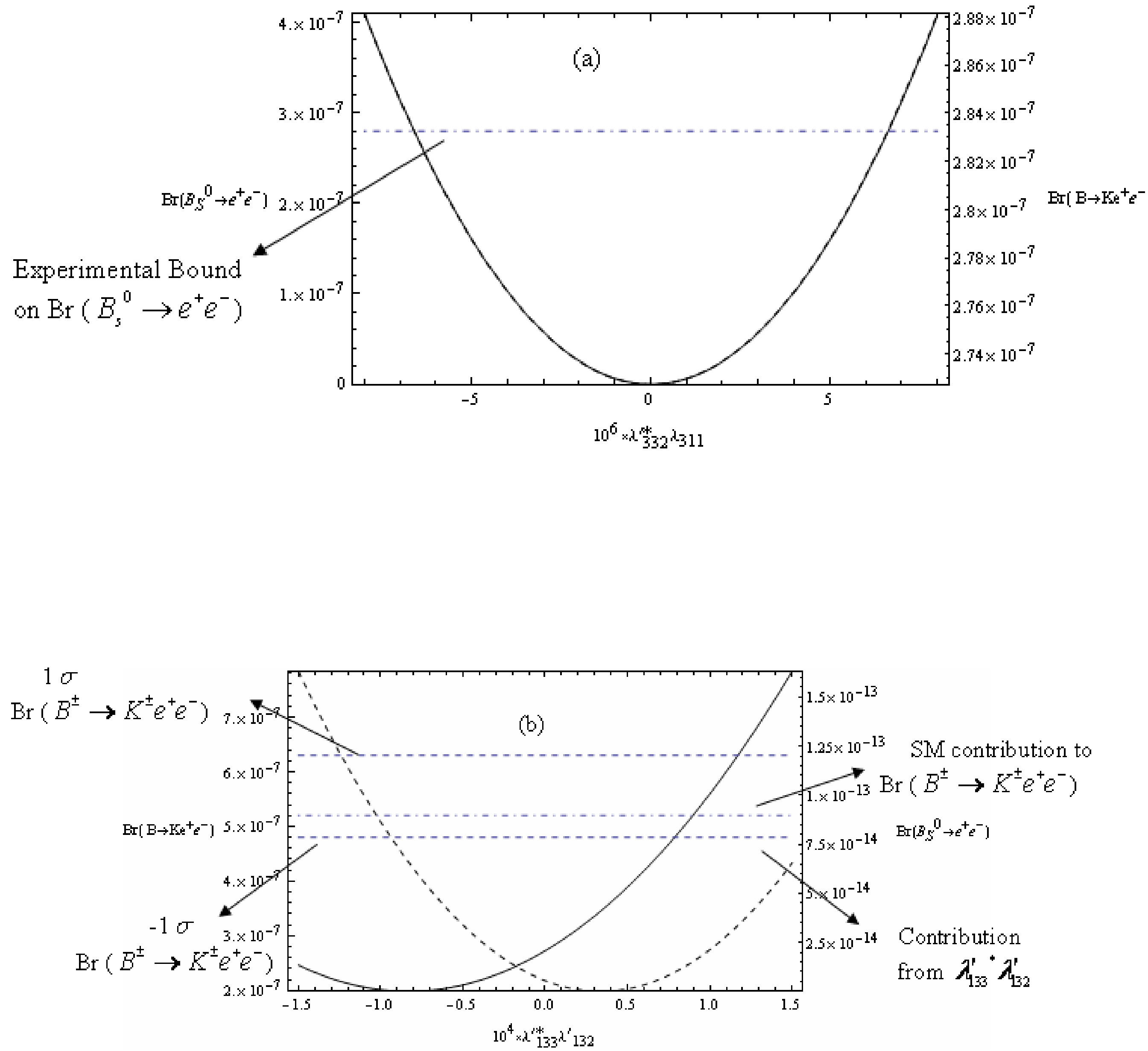


Figure 5.1.2: Comparison of branching fraction of decay process $B^\pm \rightarrow K^\pm e^+e^-$ (solid curve) with $B_s^0 \rightarrow e^+e^-$ (dashed curve) (a) as a function of sneutrino Yukawa couplings only as a function of (b) squark Yukawa couplings only. Dash-dot line (b) represents Experimental bound on $B^\pm \rightarrow K^\pm e^+e^-$. $\lambda'^*_{133}\lambda'_{132}$ is expressed in the units of $1/(m_{\widetilde{u}_{3Li_L}}/100GeV)^2$. $\lambda'^*_{332}\lambda_{311}$ is normalized to $1/(m_{\widetilde{v}_{3_L}}/100GeV)^2$.

Appendices

Appendix A

Invariant Amplitude

We reproduce the matrix element of the decay rate ($K^{\pm} \rightarrow \pi^{\pm} l^{+} l^{-}$) from[74]

$$M = C_{SS}\ \bar{s}b\ \bar{l}l + C_{VV}\ \bar{s}\gamma_{\mu}d\ \bar{l}\gamma^{\mu}l + C_{VA}\ \bar{s}\gamma_{\mu}d\ \bar{l}\gamma^{\mu}\gamma^{5}l + C_{SP}\ \bar{s}d\ \bar{l}\gamma^{5}l, \tag{A.1}$$

where, $C_{SS},\ C_{VV},\ C_{VA},\ C_{SP}$ are coefficients of interaction terms. The hadronic matrix elements are given in terms of form factors $f^{\pm}(s), f^{T}(s)$ [55, 74, 76]

$$< \pi(p_{\pi})|\bar{s}\gamma_{\mu}(1-\gamma_{5})d|K(p_{K}) >= (p_{\pi}+p_{K})_{\mu}f^{+}(s) + (p)_{\mu}f^{-}(s), \tag{A.2}$$

$$< \pi(p_{\pi})|\bar{s}b|K(p_{K}) >= \frac{(m_{K}^{2}-m_{\pi}^{2})}{m_{s}-m_{d}}f_{o}(s), \tag{A.3}$$

where $p_{\mu} = (p_{K}-p_{\pi})_{\mu}$ is the momentum transfer to dilepton pair and $s = p_{\mu}p^{\mu}$. $m_{K}, m_{\pi}, m_{d}, m_{s}$ are masses of K and π mesons, d quark and s quark respectively. We have computed the form factors $f^{\pm}(s), f^{T}(s)$ from the analytic form given by [74], which are used in our plots. $f_{o}(s)$ is defined as:

$$f_{o}(s) = f^{+}(s) + \frac{s}{m_{K}^{2}-m_{\pi}^{2}}f^{-}(s). \tag{A.4}$$

We also, have

$$M = F_{S}\ \bar{v}(l)\ u(l') + F_{V}\ p_{K_{\mu}}\ \bar{v}(l)\ \gamma^{\mu}\ u(l') + iF_{P}\ \bar{v}(l)\ \gamma^{5}\ u(l') + F_{A}\ p_{K_{\mu}}\ \bar{v}(l)\ \gamma^{\mu}\gamma^{5}\ u(l') \tag{A.5}$$

Comparing eq. (A.1) with eq. (A.5) after substituting eq. (A.2-A.4) into eq. (A.1), we get

$$\begin{aligned} F_V &= f^+(s)\ C_{VV}, \\ F_A &= f^+(s)\ C_{VA}, \\ F_S &= \frac{m_K^2}{m_s} f_o(s) C_{SS}, \\ F_P &= -i\frac{m_K^2}{m_s} f_o(s) C_{SP}, \end{aligned} \tag{A.6}$$

Appendix B

FB-Asymmetry ($K^{\pm} \to \pi^{\pm} l^{+} l^{-\prime}$)

$$M = F_S\ \bar{\text{v}}(l)\ u(l') + F_V\ p_{K_\mu}\ \bar{\text{v}}(l)\ \gamma^\mu\ u(l') + iF_P\ \bar{\text{v}}(l)\ \gamma^5\ u(l') + F_A\ p_{K_\mu}\ \bar{\text{v}}(l)\ \gamma^\mu\gamma^5\ u(l')$$

$$MM^* = (F_S\ \bar{\text{v}}(l)\ u(l') + F_V\ p_{K_\mu}\ \bar{\text{v}}(l)\ \gamma^\mu\ u(l') + iF_P\ \bar{\text{v}}(l)\ \gamma^5\ u(l') + F_A\ p_{K_\mu}\ \bar{\text{v}}(l)\ \gamma^\mu\gamma^5\ u(l')).(F_S\ \bar{\text{v}}(l)\ u(l') + F_V\ p_K^\mu\ \bar{\text{v}}(l)\ \gamma_\mu\ u(l') + iF_P\ \bar{\text{v}}(l)\ \gamma^5\ u(l') + F_A\ p_{K^\mu}\ \bar{\text{v}}(l)\ \gamma_\mu\gamma^5\ u(l'))^*$$

$$s = (P(e^+) + P(e^-))^2 = 2m_l^2 + 2EE' - 2p.p_1 = 2m_l^2 + 2P.P_1$$

$$= (P(K) - P(\pi))^2 = m_K^2 + m_\pi^2 - 2m_K E_\pi$$

$$MM^* = 4m_K^2\,|F_V|^2\,(-m_l^2 + 2EE' - \tfrac{s}{2} + m_l^2) - 4m_K^2\,|F_A|^2\,(-m_l^2 - 2EE' + \tfrac{s}{2} - m_l^2) + 4\,|F_S|^2\,(\tfrac{s}{2} - 2m_l^2) + 4\,|F_P|^2\,(p.p_1 + m_l^2) - 4m_l m_K\,\mathrm{Im}(F_A F_P^*)(m_K - E_\pi) - 8m_l m_K\,\mathrm{Re}(F_S F_V^*)(E - E')$$

$$= 2m_K^2\,|F_V|^2\,(4EE' - s) + 2m_K^2\,|F_A|^2\,(4m_l^2 + 4EE' - s) + 2\,|F_S|^2\,(s - 4m_l^2) + 2s\,|F_P|^2 - 8m_l m_K\,\mathrm{Im}(F_A F_P^*)(m_K^2 - m_\pi^2 + s) - 8m_l m_K\,\mathrm{Re}(F_S F_V^*)(E - E')$$

Three Body Phase Space

For a particular decay process $K^{\pm} \to \pi^{\pm} e^{+} e^{-}$

$$d\Gamma = \frac{(2\pi)^4 \delta^4(P_K - P_\pi - P_e - P_e')|M_A|^2 d^3P_\pi d^3P_e d^3P_e'}{(2\pi)^3(2\pi)^3(2\pi)^3 2E_\pi 2E_e' 2E_e 2m_K}$$

where P_K, P_π, P_e, P_e' are four momenta of K, π, e^-, e^+ respectively

using $\frac{d^3P_\pi}{2E_\pi} = \delta(P_\pi^2 - m_\pi^2)\ \theta(E_\pi)\ d^4P_\pi$ and Integrating over P_π

$$d\Gamma = \frac{\delta(P_\pi^2 - m_\pi^2)|M|^2 d^3P_e d^3P_e'}{256\ \pi^5 E_e' E_e m_K}$$

Replacing d^3P_e by $4\pi p_e E_e dE_e$ and $d^3P'_e$ by $2\pi p'_e E'_e dE'_e d(\cos\theta)$ and integrating over $\cos\theta$

$$d\Gamma = \frac{p_e p'_e |M|^2 \delta(m_K^2 + m_e^2 + m_e^2 - m_\pi^2 - 2m_K(E_e + E'_e) + 2P_e.P'_e) dE_e dE'_e d(\cos\theta)}{32\pi^3 m_K}$$

Integrating over $\cos\theta$, we get

$$d\Gamma = \frac{|M|^2 \delta(m_K^2 + m_e^2 + m_e^2 - m_\pi^2 - 2m_K(E_e + E'_e) + 2P_e.P'_e) dE_e dE'_e}{64\pi^3 m_K}$$

$$\frac{d^2\Gamma}{dE_e dE'_e} = \frac{1}{64\pi^3 m_K}(2m_K^2 |F_V|^2 (4EE' - s) - 2m_K^2 |F_A|^2 (-4m_e^2 - 4EE' + s) + 2|F_S|^2 (s - 4m_e^2) + 2|F_P|^2 (S) - 4m_e \operatorname{Im}(F_A F_P^*)(m_K^2 - m_\pi^2 + s) - 8m_e m_K \operatorname{Re}(F_S F_V^*)(E - E'))$$

In the K^+ rest frame,

$$s = (P_K - P_\pi)^2$$

$$= m_K^2 + m_\pi^2 - 2P_K.P_\pi$$

$$= m_K^2 + m_\pi^2 - 2m_K E_\pi$$

$$= m_K^2 + m_\pi^2 - 2m_K(m_K - E - E')$$

$$= m_K^2 + m_\pi^2 - 2m_K^2 + 2m_K E + 2m_K E'$$

$$s = m_\pi^2 - m_K^2 + 2m_K E + 2m_K E'$$

$$E + E' = \frac{s + m_K^2 - m_\pi^2}{2m_K}$$

In the e^+ e^- rest frame,

$$t = (P_K - P_e)^2 = m_K^2 + m_e^2 - 2E_K E_e + 2p_K p_e \cos\theta$$

$$t' = (P_K - P_{e'})^2 = m_K^2 + m_e^2 - 2E_K E_{e'} - 2p_K p_e \cos\theta$$

Since $E_{e'} = E_e$

$$t' = (P_K - P_{e'})^2 = m_K^2 + m_e^2 - 2E_K E_e - 2p_K p_e \cos\theta$$

$$t - t' = 4p_K p_e \cos\theta$$

In K meson rest frame

$$t = (P_K - P_e)^2 = m_K^2 + m_e^2 - 2m_K E_e$$

$$t' = (P_K - P_{e'})^2 = m_K^2 + m_e^2 - 2m_K E_{e'}$$

$$t - t' = 4m_K(E_{e'} - E_e)$$

Since $t - t'$ is invariant

$$4m_K(E_{e'} - E_e) = 4p_K p_e \cos\theta$$

$$m_K(E_{e'} - E_e) = p_K p_e \cos\theta$$

In the $e^+ \, e^-$ rest frame,

$$s = (P_e + P_{e'})^2$$

$$= 2m_e^2 + 2E^2 + 2p^2$$

$$s = 4E^2 = (\text{Total Energy of dilepton pair})^2$$

$$s = (P_K - P_e - P_{e'})^2$$

$$p_e = -p_{e'}$$

$$E_K = \frac{m_K^2 - m_\pi^2 + 4E_e^2}{4E_e}$$

$$p_K^2 = \frac{m_K^4 + m_\pi^4 + 16E_e^4 - 2m_\pi^2 m_K^2 + 8E_e^2 m_K^2 - 8E_e^2 m_\pi^2}{16E_e^2} - m_K^2$$

$$m_K(E_{e'} - E_e) = p_K p_e \cos\theta$$

$$E_{e'} - E_e = \frac{\lambda^{\frac{1}{2}}(s)}{2m_K}\sqrt{1 - \frac{4m_e^2}{s}} \cos\theta$$

$$= \frac{\lambda^{\frac{1}{2}}(s)\beta_l \cos\theta}{2m_K}$$

$$E = \frac{s+m_K^2-m_\pi^2+\lambda^{\frac{1}{2}}(s)\beta_l\cos\theta}{4m_K};\ E' = \frac{s+m_K^2-m_\pi^2-\lambda^{\frac{1}{2}}(s)\beta_l\cos\theta}{4m_K}$$

Using above relations and Jacobian Determinant $J(s,\theta)$ we can re-write

$$\frac{d^2\Gamma}{ds\,d\cos\theta} = \frac{1}{256\pi^3 m_K^3}\beta_e\lambda^{\frac{1}{2}}(s)(\tfrac{1}{4}\left|F_V\right|^2\lambda(s)(1-\beta_e^2\cos^2\theta)+\left|F_A\right|^2(\tfrac{1}{4}\lambda(s)(1-\beta_e^2\cos^2\theta)+ 4m_K^2m_e^2)+s\beta_e^2\left|F_S\right|^2+s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)2m_e(m_\pi^2-m_K^2-s)+2m_e\,\mathrm{Re}(F_SF_V^*)\beta_e\lambda^{\frac{1}{2}}(s)\cos\theta)$$

Integrating (1) w.r.t. cosθ, bounded by, $1 \leq \cos\theta \geq -1$;we get

$$= \frac{1}{256\pi^3 m_K^3}(\beta_e\lambda^{\frac{1}{2}}(s)((\tfrac{1}{4}\left|F_V\right|^2\lambda(s)(2-\tfrac{2}{3}\beta_e^2)+A^2(\tfrac{1}{4}\lambda(s)(2-\tfrac{2}{3}\beta_e^2)+8m_K^2m_e^2)+s\beta_e^2\left|F_S\right|^2+ s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)4m_e(m_\pi^2-m_K^2-s))$$

$$\beta_l = \sqrt{1-\frac{4m_e^2}{s}}$$

$$\frac{d\Gamma}{ds} = \frac{1}{256\pi^3 m_K^3}\beta_e\lambda^{\frac{1}{2}}(s)(\tfrac{1}{2}\left|F_V\right|^2\lambda(s)(1-\tfrac{1}{3}+\tfrac{4m_e^2}{3F_S})+\left|F_A\right|^2(\tfrac{1}{2}\lambda(s)(1-\tfrac{1}{3}+\tfrac{4m_e^2}{3s})+8m_K^2m_e^2)+ 2s\beta^2\left|F_S\right|^2+2s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)4m_e(m_\pi^2-m_K^2-s)))$$

$$= \frac{1}{256\pi^3 m_K^3}\beta_e\lambda^{\frac{1}{2}}(s)(\tfrac{1}{3}\left|F_V\right|^2\lambda(s)(1+\tfrac{2m_e^2}{s})+\left|F_A\right|^2(\tfrac{1}{3}\lambda(s)(1+\tfrac{2m_e^2}{s})+8m_K^2m_e^2)+s\beta_e^2\left|F_S\right|^2+ s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)4m_e(m_\pi^2-m_K^2-s))$$

$$A_{FB}(s) = \frac{\int_0^1 d(\cos\theta)\frac{d^2\Gamma}{ds\,d(\cos\theta)} - \int_{-1}^0 d(\cos\theta)\frac{d^2\Gamma}{ds\,d(\cos\theta)}}{\int_0^1 d(\cos\theta)\frac{d^2\Gamma}{ds\,d(\cos\theta)} + \int_{-1}^0 d(\cos\theta)\frac{d^2\Gamma}{ds\,d(\cos\theta)}} \tag{B.1}$$

$$\frac{d^2\Gamma}{ds\,d(\cos\theta)} = \frac{1}{256\pi^3 m_K^3}\beta_l\lambda^{\frac{1}{2}}(s)(\tfrac{1}{4}\left|F_V\right|^2\lambda(s)(1-\beta_l^2\cos^2\theta)+\left|F_A\right|^2(\tfrac{1}{4}\lambda(s)(1-\beta_l^2\cos^2\theta)+ 4m_K^2m_l^2)+_S\beta_l^2\left|F_S\right|^2+s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)2m_l(m_\pi^2-m_K^2-s)+2m_l\,\mathrm{Re}(F_SF_V^*)\beta_l\lambda^{\frac{1}{2}}(s)\cos\theta)$$

Integrating w.r.t. cosθ, bounded by, $1 \geq \cos\theta \geq 0$;we get

$$\begin{aligned}\frac{d\Gamma}{ds}|_0^1 &= \frac{1}{256\pi^3 m_K^3}\beta_e\lambda^{\frac{1}{2}}(s)(\frac{1}{4}\left|F_V\right|^2\lambda(s)(1-\frac{1}{3}\beta_e^2)+ \\ &\quad \left|F_A\right|^2(\frac{1}{4}\lambda(s)(1-\frac{1}{3}\beta_e^2)+4m_K^2m_e^2)+ \\ &\quad s\beta_l^2\left|F_S\right|^2+s\left|F_P\right|^2+\mathrm{Im}(F_AF_P^*)2m_e(m_\pi^2-m_K^2-s)+m_e\,\mathrm{Re}(F_SF_V^*)\beta_e\lambda^{\frac{1}{2}}(s))\end{aligned} \tag{B.2}$$

Integrating w.r.t. cosθ, bounded by, $-1 \leq \cos\theta \leq 0$;we get

$$\frac{d\Gamma}{ds}|_{-1}^{0} = \frac{1}{256\pi^3 m_K^3}\beta_e \lambda^{\frac{1}{2}}(s)(\frac{1}{4}\left|F_V\right|^2 \lambda(s)(1-\frac{1}{3}\beta_e^2) + \left|F_A\right|^2(\frac{1}{4}\lambda(s)(1-\frac{1}{3}\beta_e^2) + 4m_K^2 m_l^2) + s\beta_e^2\left|F_S\right|^2 + s\left|F_P\right|^2 + \mathrm{Im}(F_A F_P^*)2m_e(m_\pi^2 - m_K^2 - s) - m_e \,\mathrm{Re}(F_S F_V^*)\beta_e \lambda^{\frac{1}{2}}(s)) \tag{B.3}$$

$$\frac{d\Gamma}{ds} = \frac{d\Gamma}{ds}|_0^1 + \frac{d\Gamma}{ds}|_{-1}^0$$

Adding eq. (B.2&B.3), we get [46]

$$\frac{d\Gamma}{ds} = \frac{1}{256\pi^3 m_K^3}\beta_e \lambda^{\frac{1}{2}}(s)(\frac{1}{2}\left|F_V\right|^2 \lambda(s)(1-\frac{1}{3}\beta_e^2) + \left|F_A\right|^2(\frac{1}{2}\lambda(s)(1-\frac{1}{3}\beta_e^2) + 8m_K^2 m_e^2) + s\beta_e^2\left|F_S\right|^2 + s\left|F_P\right|^2 + 4m_l(m_\pi^2 - m_K^2 - s)\,\mathrm{Im}(F_A F_P^*)) \tag{B.4}$$

$$\frac{d\Gamma}{ds}|_0^1 - \frac{d\Gamma}{ds}|_{-1}^0 = \frac{1}{256\pi^3 m_K^3} 2m_e \,\mathrm{Re}(F_S F_V^*)\beta_e^2 \lambda(s) \tag{B.5}$$

Using eq. (B.4 and B.5) in eq. (B.1), we get

$$A_{FB}(s) = \frac{2m_e \,\mathrm{Re}(F_S F_V^*)\beta_e \lambda^{\frac{1}{2}}(s)}{\frac{1}{2}|F_V|^2\lambda(s)(1-\frac{1}{3}\beta_e^2) + |F_A|^2(\frac{1}{2}\lambda(s)(1-\frac{1}{3}\beta_e^2) + 8m_K^2 m_e^2) + s\beta_e^2|F_S|^2 + s|F_P|^2 + 4m_e(m_\pi^2 - m_K^2 - s)\,\mathrm{Im}(F_A F_P^*)}$$

$$A_{FB}(s) = \frac{1}{256\pi^3 m_K^3} 2m_e \,\mathrm{Re}(F_S F_V^*)\beta_e^2 \lambda(s)(\frac{d\Gamma}{ds})^{-1} \tag{B.6}$$

Appendix C

Decay rate and A_{LP} of $B_q \to l_\beta l'_\beta$

$$M = C_{PP}\ (\bar{q}\gamma^5 b)(\overline{l_\beta}\gamma^5 l'_\beta) + C_{PS}\ (\bar{q}\gamma^5 b)(\overline{l_\beta} l'_\beta) + C_{AA}\ (\bar{q}\gamma^\mu\gamma^5 b)(\overline{l_\beta}\gamma_\mu\gamma^5 l'_\beta) \tag{C.1}$$

Using $PCAC$

$$< 0|\bar{q}\gamma_\mu\gamma_5 b|B >= -if_{B_q}(p_B)_\mu \tag{C.2}$$

where $p_{B_\mu} = (P_l + P_{l'})_\mu$

$$< 0|\bar{q}\gamma_5 b|B >= if_{B_q}\frac{m^2_{B_q}}{m_b + m_q} \tag{C.3}$$

Substituting eqs. (C.2 and C.3) in eq. (C.1) we get

$$M = if_{B_q}\frac{m^2_{B_q}}{m_b + m_q}C_{PP}\ (\overline{l_\beta}\gamma^5 l'_\beta) + if_{B_q}\frac{m^2_{B_q}}{m_b + m_q}C_{PS}\ (\overline{l_\beta} l'_\beta) - if_{B_q}C_{AA}\ (\overline{l_\beta}\gamma_\mu\gamma^5 l'_\beta) \tag{C.4}$$

$$M = F_S\ \bar{u}(l)\mathrm{v}(l') + F_P\ \bar{u}(l)\ \gamma^5\ \mathrm{v}(l') + F_A\ p_{B_\mu}\ \bar{u}(l)\gamma^\mu\gamma^5\mathrm{v}(l') \tag{C.5}$$

On comparing eq. (C.5) with eq. (C.4), we get

$$F_S = if_{B_q}\frac{m_{B_q}^2}{m_b+m_q}C_{PP};\; F_A=-if_{B_q} \tag{C.6}$$
$$F_P = if_{B_q}\frac{m_{B_q}^2}{m_b+m_q}C_{PS}.$$

The decay rate for $B_q\rightarrow l_\beta l'_\beta$ is given by

$$MM^* = (F_S\ \bar{u}(l)\mathrm{v}(l') + F_V\ p_{B_\mu}\ \bar{u}(l)\ \gamma^\mu \mathrm{v}(l') + F_P\ \bar{u}(l)\ \gamma^5\ \mathrm{v}(l') + F_A\ p_{B_\mu}\ \bar{u}(l)\gamma^\mu\gamma^5\mathrm{v}(l'))$$
$$(F_S\ \bar{u}(l)v(l') + F_V\ p_{B_\mu}\ \bar{u}(l)\ \gamma^\mu v(l') + F_P\ \bar{u}(l)\ \gamma^5\ v(l') + F_A\ p_{B_\mu}\ \bar{u}(l)\gamma^\mu\gamma^5 v(l'))^*$$

$$= 4\,|F_S|^2\,(p.p_1 - m^2) + 4\,|F_P + 2m_l F_A|^2\,(p.p_1 + m^2)$$

In the B^0 rest frame

$$p.p_1 = E_l^2 + p^2$$

$$|M|^2 = 8P^2\,|F_S + 2m_l F_V|^2 + 8E_l^2\,|F_P + 2m_l F_A|^2$$

Two Body Phase Spcae

$$d\Gamma = \frac{(2\pi)^4\delta^4(P_B - P_l - P_{l'})|M|^2 d^3P_l d^3P_{l'}}{(2\pi)^3(2\pi)^3 2E_{l'}2E_l 2m_B}$$

where P_B, P_l, P_l' are four momenta of B, l^-, l^+ respectively

$$d\Gamma = \frac{\delta^4(P_B - P_l - P_{l'})|M|^2 d^3P_l d^3P_{l'}}{32\pi^2 E_{l'}E_l m_B}$$

$$= \frac{\delta(m_B - E_l - E_{l'})\delta^3(p_{l'} + p_l)|M|^2 d^3P_l d^3P_{l'}}{32\pi^2 E_{l'}E_l m_B}$$

$$= \frac{\delta(m_B - E_l - E_{l'})|M|^2 d^3P_l}{32\pi^2 E_{l'}E_l m_B}$$

Since there is no angular dependence

$$= \frac{\delta(m_B - E_l - E_{l'})|M|^2 p_l dE_l}{8\pi E_{l'} m_B}$$

since

$$E_{l'} = E_l \text{(Lepton Flavor conserving case)}$$
$$= \frac{\delta(m_B - 2E_l)|M|^2 p_l dE_l}{8\pi E_l m_B}$$

$$\text{Using } p_l = E_l(1 - \frac{m_l^2}{E_l^2})^{\frac{1}{2}}$$

$$\Gamma = \frac{|M|^2(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}}}{16\pi m_B}$$

$$|M|^2 = 8P^2 |F_S|^2 + 8E_l^2 |F_P + 2m_l F_A|^2$$

$$= 8E_l^2(1 - \frac{m_l^2}{E_l^2}) |F_S|^2 + 8E_l^2 |F_P + 2m_l F_A|^2$$

$$= 8E_l^2[(1 - \frac{m_l^2}{E_l^2}) |F_S|^2 + |F_P + 2m_l F_A|^2]$$

$$= \frac{|M|^2}{8\pi m_B}(1 - \frac{4m_l^2}{m_B^2})^{\frac{1}{2}}$$

$$\Gamma = \frac{m_B}{8\pi}(1 - \frac{4m_l^2}{m_B^2})^{\frac{1}{2}}[(1 - \frac{4m_l^2}{m_B^2}) |F_S|^2 + |F_P + 2m_l F_A|^2] \tag{C.7}$$

We write the l^+ four spin vector in terms of a unit vector, ξ, along the l^+ spin in its rest frame, as

$$s_+^0 = \frac{\vec{P}_+ \widehat{\xi}}{m_l}; \quad \vec{s}_+ = \widehat{\xi} + \frac{s_+^0}{E_l + m_l}\vec{p}_+$$
$$\widehat{s}_l = \frac{\vec{p}_+}{|\vec{p}_+|};$$

The longitudinal polarization asymmetry is calculated as

$$A_{LP}^{\pm} = \frac{[\Gamma(s_{l^-}, s_{l^+}) + \Gamma(\mp s_{l^-}, \pm s_{l^+})] - [\Gamma(\pm s_{l^-}, \mp s_{l^+}) + \Gamma(-s_{l^-}, -s_{l^+})]}{[\Gamma(s_{l^-}, s_{l^+}) + \Gamma(\mp s_{l^-}, \pm s_{l^+})] + [\Gamma(\pm s_{l^-}, \mp s_{l^+}) + \Gamma(-s_{l^-}, -s_{l^+})]} \tag{C.8}$$

Left and Right handed helicity spinors for particles/antiparticles are [83]:

$$u\uparrow = \sqrt{E+m}\begin{pmatrix} \cos\frac{\theta}{2} \\ e^{i\phi}\sin\frac{\theta}{2} \\ \frac{|\vec{p}|}{E+m}\cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E+m}e^{i\phi}\sin\frac{\theta}{2} \end{pmatrix};\ u\downarrow = \sqrt{E+m}\begin{pmatrix} -\sin\frac{\theta}{2} \\ e^{i\phi}\cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E+m}\sin\frac{\theta}{2} \\ -\frac{|\vec{p}|}{E+m}e^{i\phi}\cos\frac{\theta}{2} \end{pmatrix};$$

$$\text{v}\uparrow = \sqrt{E+m}\begin{pmatrix} \frac{|\vec{p}|}{E+m}\sin\frac{\theta}{2} \\ -\frac{|\vec{p}|}{E+m}e^{i\phi}\cos\frac{\theta}{2} \\ -\sin\frac{\theta}{2} \\ e^{i\phi}\cos\frac{\theta}{2} \end{pmatrix};\ \text{v}\downarrow = \sqrt{E+m}\begin{pmatrix} \frac{|\vec{p}|}{E+m}\cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E+m}e^{i\phi}\sin\frac{\theta}{2} \\ \cos\frac{\theta}{2} \\ e^{i\phi}\sin\frac{\theta}{2} \end{pmatrix}$$

For the final state l^+ which has polar angle θ and $\phi = 0$

$$u\uparrow = \sqrt{E_l+m}\begin{pmatrix} \cos\frac{\theta}{2} \\ \sin\frac{\theta}{2} \\ \frac{|\vec{p}|}{E_l+m}\cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E_l+m}\sin\frac{\theta}{2} \end{pmatrix};\ u\downarrow = \sqrt{E_l+m}\begin{pmatrix} -\sin\frac{\theta}{2} \\ \cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E_l+m}\sin\frac{\theta}{2} \\ -\frac{|\vec{p}|}{E_l+m}\cos\frac{\theta}{2} \end{pmatrix};$$

For the final state l^- which has polar angle $\pi - \theta$ and $\phi = \pi$

$$\text{v}\uparrow = \sqrt{E_l+m}\begin{pmatrix} \frac{|\vec{p}|}{E_l+m}\cos\frac{\theta}{2} \\ \frac{|\vec{p}|}{E_l+m}\sin\frac{\theta}{2} \\ -\cos\frac{\theta}{2} \\ -\sin\frac{\theta}{2} \end{pmatrix};\ \text{v}\downarrow = \sqrt{E_l+m}\begin{pmatrix} \frac{|\vec{p}|}{E_l+m}\sin\frac{\theta}{2} \\ -\frac{|\vec{p}|}{E_l+m}\cos\frac{\theta}{2} \\ \sin\frac{\theta}{2} \\ -\cos\frac{\theta}{2} \end{pmatrix}$$

$$M(\uparrow\uparrow) = (F_S)\ \bar{u}\uparrow(l)\ \ \text{v}\uparrow(l') + (F_P + 2m_eF_A)\bar{u}\uparrow(l)\ \gamma^5\ \text{v}\uparrow(l') = \ 2p(F_P + 2m_eF_A) - 2E_lF_S$$

$$M(\downarrow\downarrow) = (F_S)\ \bar{u}\downarrow(l)\ \ \text{v}\downarrow(l') + (F_P + 2m_eF_A)\bar{u}\downarrow(l)\ \gamma^5\ \text{v}\downarrow(l') = -2p(F_P + 2m_eF_A) - 2E_lF_S$$

$$M(\uparrow\downarrow) = M(\downarrow\uparrow) = 0$$

$$\Gamma(\uparrow\uparrow) = \frac{|M(\uparrow\uparrow)|^2(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}}}{32\pi m_B}$$

$$|M(\uparrow\uparrow)|^2 = 4(P^2\,|F_S|^2 + E_l^2\,|F_P + 2m_lF_A|^2)$$

$$\Gamma(\uparrow\uparrow) = \frac{|M(\uparrow\uparrow)|^2(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}}}{32\pi m_B}$$

$$= \frac{(m_B)^2((1-\frac{4(m_l)^2}{(m_B)^2})|F_S|^2+(m_B)^2|F_P+2m_lF_A|^2)(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}} - \mathrm{Re}[\sqrt{1-\frac{4(m_l)^2}{(m_B)^2}}F_S(F_P+2m_lF_A)^*])}{8\pi m_B}$$

$$\Gamma(\uparrow\uparrow) = \frac{m_B((1-\frac{4(m_l)^2}{(m_B)^2})\,|F_S|^2 + |F_P+2m_lF_A|^2)(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}} - \mathrm{Re}[\sqrt{1-\frac{4(m_l)^2}{(m_B)^2}}F_S(F_P+2m_lF_A)^*])}{8\pi} \tag{C.9}$$

Similarly

$$\Gamma(\downarrow\downarrow) = \frac{m_B((1-\frac{4(m_l)^2}{(m_B)^2})\,|F_S|^2 + |F_P+2m_lF_A|^2)(1-\frac{4m_l^2}{m_B^2})^{\frac{1}{2}}) + \mathrm{Re}[\sqrt{1-\frac{4(m_l)^2}{(m_B)^2}}F_S(F_P+2m_lF_A)^*])}{8\pi} \tag{C.10}$$

eq. (C.8) becomes

$$A_{LP} = \frac{[\Gamma(\uparrow\uparrow)+\Gamma(\uparrow\downarrow)] - [\Gamma(\downarrow\downarrow)+\Gamma(\downarrow\uparrow)]}{[\Gamma(\uparrow\uparrow)+\Gamma(\uparrow\downarrow)] + [\Gamma(\downarrow\downarrow)+\Gamma(\downarrow\uparrow)]}$$

$$A_{LP} = \frac{\Gamma(\uparrow\uparrow)-\Gamma(\downarrow\downarrow)}{\Gamma(\uparrow\uparrow)+\Gamma(\downarrow\downarrow)} \tag{C.11}$$

Since $\Gamma(\uparrow\downarrow) = \Gamma(\downarrow\uparrow) = 0$ (conservation of angular momentum).Substituting eq. (C.9 and C.10) in eq. (C.11) we get

$$A_{LP} = -2\sqrt{1-\frac{4(m_l)^2}{(m_{B_q})^2}}\frac{\mathrm{Re}[F_S(F_P+2m_lF_A)^*]}{[(1-\frac{4m_l^2}{(m_{B_q})^2})\,|F_S|^2 + |F_P+2m_lF_A|^2]} \tag{C.12}$$

After Substiting from eq. (C.6)

$$A_{LP} = \frac{2\frac{(m_{B_q})}{m_b+m_q}\sqrt{1-\frac{4(m_l)^2}{(m_{B_q})^2}}\,\mathrm{Re}[C_{PS}(2\frac{m_l}{m_{B_q}}C_{AA} - \frac{(m_{B_q})}{m_b+m_q}C_{PP})^*]}{|\,\frac{2m_l}{m_{B_q}}C_{AA} - \frac{(m_{B_q})}{m_b+m_q}C_{PP}\,|^2 + (1-\frac{4(m_l)^2}{(m_{B_q})^2})\left|\frac{(m_{B_q})}{m_b+m_q}C_{PS}\right|^2} \tag{C.13}$$

Appendix D

Wilson Coefficients

We reproduce the matrix element of the decay rate ($B \to K l^+ l^-$) from [74]

$$M = F_S \bar{l} l + F_V (p_B)_\mu \bar{l} \gamma^\mu l + F_A (p_B)_\mu \bar{l} \gamma^\mu \gamma^5 l + F_P \bar{l} \gamma^5 l$$

The hadronic matrix elements contained in $F_{V,A,S,P}$ are given in terms of form factors $f^{\pm}(s), f^T(s)$ [55, 74, 76]

$$< K(p_K)|\bar{s}\gamma_\mu(1-\gamma_5)b|B(p_B) >= (p_B + p_K)_\mu f^+(s) + p_\mu f^-(s) \tag{D.1}$$

$$< K(p_K)|\bar{s} i \sigma_{\mu v}\gamma^v(1+\gamma_5)b|B(p_B) >= (p_B + p_K)_\mu s - p_\mu(m_B^2 - m_K^2)\frac{f^T(s)}{m_B + m_K},$$

$$< K(p_K)|\bar{s}b|B(p_B) >= \frac{(m_B^2 - m_K^2)}{m_b - m_s} f_o(s), \tag{D.2}$$

where $p_\mu = (p_B - p_K)_\mu$ is the momentum transfer to dilepton pair and $s = p_\mu p^\mu$. m_B, m_K, m_b, m_s are masses of B and K mesons, b quark and s quark respectively. We have computed form factors $f^{\pm}(s), f^T(s)$ from the analytic form given by [74], which are used in our plots. $f_o(s)$ is defined as

$$f_o(s) = f^+(s) + \frac{s}{m_B^2 - m_K^2} f^-(s) \tag{D.3}$$

The effective Hamiltonian for the given decay process is given by [74]

$$H_{eff} = \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^{10} \{C_i(\mu)\ O_i(\mu) + C_{Q_i}(\mu)\ O_{Q_i}(\mu)\}. \tag{D.4}$$

C_1	C_2	C_3	C_4	C_5	C_6
-0.249	1.107	0.011	-0.025	0.007	-0.031

Table 5.1.2: A table of numerical values of Wilson coefficients[54].

The Wilson coefficients relevant to our calculation are evaluated at $\mu = m_b$:

$$C_7 = -0.315,\ C_{10} = -4.642. \tag{D.5}$$

The expression for C_9^{eff} in the next to leading order approximation used here is given by [55, 74, 76]

$$C_9^{eff}(\widehat{m}_b, \widehat{s}) = A(\widehat{s}) + B(\widehat{s}) \tag{D.6}$$

where

$$A(\widehat{s}) = 4.227 + 0.124w(\widehat{s}) + \sum_{i=1}^{6} \alpha_i (\widehat{s}) C_i \tag{D.7}$$

$$\alpha_1 (\widehat{s}) = 3g(\widehat{m}_c, \widehat{s});\ \alpha_2 (\widehat{s}) = g(\widehat{m}_c, \widehat{s});\ \alpha_3 (\widehat{s}) = 3g(\widehat{m}_c, \widehat{s}) - \tfrac{1}{2}g(\widehat{m}_s, \widehat{s}) - 2g(\widehat{m}_b, \widehat{s}) + \tfrac{2}{3};$$
$$\alpha_4 (\widehat{s}) = \tfrac{2}{9} + g(\widehat{m}_c, \widehat{s}) - \tfrac{3}{2}g(\widehat{m}_s, \widehat{s}) - 2g(\widehat{m}_b, \widehat{s});$$
$$\alpha_5 (\widehat{s}) = \tfrac{2}{3} + 3g(\widehat{m}_c, \widehat{s}) - \tfrac{3}{2}g(\widehat{m}_b, \widehat{s});\ \alpha_6 (\widehat{s}) = \tfrac{2}{9} + g(\widehat{m}_c, \widehat{s}) - \tfrac{1}{2}g(\widehat{m}_b, \widehat{s})$$

$$B(\widehat{s}) = \lambda_t(3C_1 + C_2)(g(\widehat{m}_c, \widehat{s}) - g(\widehat{m}_u, \widehat{s})) \tag{D.8}$$

where coefficients $C_1 - C_6$(evaluated at m_b), $g(\widehat{m}_q, \widehat{s})$ and $w(\widehat{s})$ [55, 74, 76] used in our computation are listed as:

$$\begin{aligned} g(\widehat{m}_q, \widehat{s}) \;=\; & \frac{-8}{9}\log[\widehat{m}_q] + \frac{4}{9}y_q - \frac{2}{9}(2 + y_q)\sqrt{1 - y_q} + \\ & \{\theta(1 - y_q)(\log(\frac{1 + \sqrt{1 - y_q}}{1 - \sqrt{1 - y_q}}) - i\pi) + \theta(y_q - 1)\tan^{-1}(\frac{1}{\sqrt{y_q - 1}})\} \end{aligned} \tag{D.9}$$

$$w(s) = -\frac{2}{9}\pi^2 - \frac{4}{3}Li_2(\widehat{s}) - \frac{2}{3}\ln[\widehat{s}]\ln(1\text{-}\widehat{s}) - \frac{5+4\widehat{s}}{3(1+2\widehat{s})}\ln(1-\widehat{s}) - \frac{2\widehat{s}(1+\widehat{s})(1-2\widehat{s})}{3(1-\widehat{s})^2(1+2\widehat{s})}\ln[\widehat{s}] + \frac{5+9\widehat{s}-6(\widehat{s})^2}{6(1-\widehat{s})(1+2\widehat{s})} \tag{D.10}$$

where

$\lambda_t = \frac{V_{ub}^* V_{us}}{V_{tb}^* V_{ts}}$, $\widehat{m}_q = \frac{m_q}{m_b}$, $\widehat{s} = \frac{s}{(m_b)^2}$, $y_q = \frac{(2\widehat{m}_q)^2}{\widehat{s}}$

C_7 belongs to photon Penguin and C_9^{eff} and C_{10} belong to W box and Z Penguin Feynman diagrams contributing to the process under discussion here. B is continuum part of $u\bar{u}$ and $c\bar{c}$ loops proportional to $V_{ub}^* V_{uq}$ and $V_{cb}^* V_{cq}$ respectively.

Appendix E

CP-Asymmetry

We calculate the A_{CP} as

$$A_{CP} = \frac{\frac{d\Gamma(B\to Kl^+l^-)}{ds} - \frac{d\Gamma(\bar{B}\to\bar{K}l^-l^+)}{ds}}{\frac{d\Gamma(B\to Kl^+l^-)}{ds} + \frac{d\Gamma(\bar{B}\to\bar{K}l^-l^+)}{ds}} \tag{E.1}$$

Since

$$\begin{aligned}\frac{d\Gamma(\bar{B}\to\bar{K}l^-l^+)}{ds} &= \frac{1}{256\pi^3 m_B^3}\beta(m_l,s)\lambda^{\frac{1}{2}}(s)\{\left|\bar{F}_S\right|^2 2s\beta_l^2 + \left|\bar{F}_P\right|^2 2s \\ &+ \left|\bar{F}_V\right|^2 \frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) + \left|\bar{F}_A\right|^2 [\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) \\ &+8m_B^2 m_l^2] + \mathrm{Re}(\bar{F}_P^*\bar{F}_A)4m_l(m_B^2 - m_K^2 + s)\},\end{aligned} \tag{E.2}$$

$$\begin{aligned}\frac{d\Gamma(B\to Kl^+l^-)}{ds} &= \frac{1}{256\pi^3 m_B^3}\beta(m_l,s)\lambda^{\frac{1}{2}}(s)\{|F_S|^2 2s\beta_l^2 + |F_P|^2 2s \\ &+ |F_V|^2 \frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) + |F_A|^2 [\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) \\ &+8m_B^2 m_l^2] + \mathrm{Re}(F_P^* F_A)4m_l(m_B^2 - m_K^2 + s)\},\end{aligned} \tag{E.3}$$

Substituting eq. (E.2 and E.3) in eq. (E.1), we get

$$A_{CP} = \lambda(s)(1+\frac{2m_l^2}{s})(|F_V|^2 - \left|\bar{F}_V\right|^2)(D(s))^{-1} \tag{E.4}$$

where

$$D(s) = 6(|F_S|^2\, 2s\beta(m_l,s)^2 + |F_P|^2\, 2s + |F_A|^2\, [\frac{1}{3}\lambda(s)(1+\frac{2m_l^2}{s}) + 8m_B^2 m] + \mathrm{Re}(F_P F_A^*)4m_l(m_B^2 - m_K^2 + s) + \frac{\lambda(s)}{6}(1+\frac{2m_l^2}{s})(|F_V|^2 + |\bar{F}_V|^2))$$

And

$$\begin{aligned}
F_V &= \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}(2C_9^{eff}(\widehat{m}_b,\widehat{s})f^+(s) - C_7\frac{4m_b}{m_B+m_K}f^T(s)) + \frac{1}{4}f^+(s)\sum_{i,m,n=1}^{3} V_{ni}^\dagger V_{im}\frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{m^2_{\widetilde{u_i^c}}},\\
F_A &= \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}(2C_{10}f^+(s)) - \frac{1}{4}f^+(s)\sum_{i,m,n=1}^{3} V_{ni}^\dagger V_{im}\frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{m^2_{\widetilde{u_i^c}}},\\
F_S &= -\frac{1}{2}\frac{(m_B^2-m_K^2)}{m_b-m_s}f_o(s)\sum_{i=1}^{3}\frac{(\lambda^*_{i\beta\beta}\lambda'_{i23} + \lambda_{i\beta\beta}\lambda'^*_{i32})}{m^2_{\widetilde{\nu}_{Li}}}, \qquad (E.5)\\
F_P &= \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}2m_l C_{10}(f^+(s)+f^-(s)) + \frac{1}{2}\frac{(m_B^2-m_K^2)}{m_b-m_s}f_o(s)\sum_{i=1}^{3}\frac{(\lambda_{i\beta\beta}\lambda'^*_{i32} - \lambda^*_{i\beta\beta}\lambda'_{i23})}{m^2_{\widetilde{\nu}_{Li}}},
\end{aligned}$$

$$\begin{aligned}
\bar{F}_V &= \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}(2\bar{C}_9^{eff}f^+(s) - C_7\frac{4m_b}{m_B+m_K}f^T(s)) + \frac{1}{4}f^+(s)(\sum_{m,n,i=1}^{3} V_{ni}^\dagger V_{im}\frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{m^2_{\widetilde{u_i^c}}})^*,\\
\bar{F}_A &= F_A^*,\ \bar{F}_S = F_S^*,\ \bar{F}_P = F_p^*, \bar{C}_9^{eff}(\widehat{m}_b,\widehat{s}) = C_9^{eff}(\widehat{m}_b,\widehat{s},\lambda_t^*), \qquad (E.6)\\
\lambda'_{\beta n3}\lambda'^*_{\beta m2} &= \left|\lambda'_{\beta n3}\lambda'^*_{\beta m2}\right| e^{i\Theta}, \beta = e,\mu
\end{aligned}$$

eq. (4.2.19) for A_{CP} contains the following factor

$$\begin{aligned}
|F_V|^2 - |\bar{F}_V|^2 &= 4(\frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi})^2 f^+(s)\,\mathrm{Im}(A'(s)B'(s)^*)\,\mathrm{Im}(\lambda_t) +\\
&\quad \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}f^+(s)\,\mathrm{Im}(A'(s))\,\mathrm{Im}(\sum_{m,n,i=1}^{3} V_{ni}^\dagger V_{im}\frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{2m^2_{\widetilde{u_i^c}}}) + \qquad (E.7)\\
&\quad \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}(f^+(s))^2\,\mathrm{Im}(B'(s))\,\mathrm{Im}(\lambda_t^*\sum_{m,n,i=1}^{3} V_{ni}^\dagger V_{im}\frac{\lambda'_{\beta n3}\lambda'^*_{\beta m2}}{2m^2_{\widetilde{u_i^c}}}),
\end{aligned}$$

where

$$A'(s) = 2Af^+(s) - C_7 \frac{4m_b}{m_B + m_K} f^T(s)$$
$$B'(s) = 2(3C_1 + C_2)(g(\widehat{m}_c, \widehat{s}) - g(\widehat{m}_u, \widehat{s})),$$

$$\begin{aligned}
|F_V|^2 + |\bar{F}_V|^2 &= G(s) + H(s), \\
G(s) &= 2(f^+(s))^2 (\left| \sum_{m,n,i=1}^{3} V_{ni}^\dagger V_{im} \frac{\lambda'_{\beta n3} \lambda'^*_{\beta m2}}{m^2_{\widetilde{u^c_i}}} \right|)^2 + \\
&\quad 4(\frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi}) f^+(s) \,\mathrm{Re}(A') \,\mathrm{Re}(\sum_{m,n,i=1}^{3} V_{ni}^\dagger V_{im} \frac{\lambda'_{\beta n3} \lambda'^*_{\beta m2}}{2m^2_{\widetilde{u^c_i}}}) + \\
&\quad 4\frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi} (f^+(s))^2 \,\mathrm{Re}(\lambda_t^* \sum_{i,m,n=1}^{3} V_{ni}^\dagger V_{im} \frac{\lambda'_{\beta n3} \lambda'^*_{\beta m2}}{2m^2_{\widetilde{u^c_i}}}) \,\mathrm{Re}(B'), \\
H(s) &= (\frac{G_F \alpha V_{tb} V_{ts}^*}{2\sqrt{2}\pi})^2 (\left| 2C_9^{eff}(\widehat{m}_b, \widehat{s}) f^+(s) - C_7 \frac{4m_b}{m_B + m_K} f^T(s) \right|^2 + \\
&\quad \left| 2\bar{C}_9^{eff}(\widehat{m}_b, \widehat{s}) f^+(s) - C_7 \frac{4m_b}{m_B + m_K} f^T(s) \right|^2)
\end{aligned} \tag{E.8}$$

The factors contributing to the CP- asymmetry at various energies are;

$$\begin{aligned}
\mathrm{Im}(A'(s)) &= 2f^+(s) \,\mathrm{Im}(A(s) \\
&= 2f^+(s) \left| \begin{array}{c} \frac{1}{2}\pi(C_3 + 3C_4) \\ -\pi(3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) \\ +\frac{1}{2}\pi(C_3 + 3C_4) \\ -\pi(3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) + \frac{1}{2}\pi(C_3 + 3C_4) \\ +\frac{1}{2}\pi(4C_3 + 4C_4 + 3C_5) \end{array} \right\} \begin{array}{c} m_c^2 > s > m_s^2 \\ m_b^2 > s > m_c^2 \\ \\ s > m_b^2 \end{array} \\
\mathrm{Im}(B'(s)) &= 2(3C_1 + C_2) \,\mathrm{Im}(g(\widehat{m}_c, \widehat{s}) - g(\widehat{m}_u, \widehat{s})), \\
&= 2(3C_1 + C_2) \left| \begin{array}{c} \frac{2\pi}{9}(2 + y_u)\sqrt{1 - y_u} \\ \frac{2\pi}{9}((2 + y_u)\sqrt{1 - y_u} - (2 + y_c)\sqrt{1 - y_c}) \end{array} \right\} \begin{array}{c} m_c^2 > s > m_u^2 \\ s > m_c^2 \end{array}
\end{aligned} \tag{E.9}$$

$$\lambda_t = -C\lambda^2 e^{i\delta}$$

Since for simplicity, we have considered only the $i = 3$ contribution in the sum (see eq. $(E.8)$)

for $\not{R}_p$ Yukawa couplings, V_{tb} contributes in $F_{A,V}$ only. Further, since $(V_{td}, V_{ts}) \sim 10^{-3}$, we neglect these numbers and take $V_{tb} \sim 1$ [67]. Since [77, 67]

$$0.23 < C < 0.59;\ 0.216 < \lambda < 0.223;\ \delta = (77^{+30}_{-32})^o;$$

We use the central value of C and λ_t i.e., 0.41 and 0.22 respectively. This parameterises the CP-violating phase used in Fig. (4.3.7).

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