PHENOMENOLOGICAL STUDIES OF THE OBSERVED ANOMALIES IN THE τ SECTOR.

By

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A thesis submitted to the

Board of Studies in Physical Sciences

In partial fulfillment of requirements

For the Degree of

DOCTOR OF PHILOSOPHY

of

HOMI BHABHA NATIONAL INSTITUTE

January 23 , 2018

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Declaration

I, hereby declare that the investigation presented in this thesis has been carried out by me under the guidance of my supervisor Prof. Nita Sinha. The work is original and has not been submitted earlier as a whole or in part of a degree/diploma at this or any other Institution/University.

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List of Publications arising from the thesis

1. Published:

   (a) H. Zeen Devi, L. Dhargyal and Nita Sinha.
   Can the observed CP asymmetry in $A_{CP}(\tau \to K_S \pi \nu)$ be due to non-standard tensor interactions? [Phys. Rev. D90, 013016]

   (b) Lobsang Dhargyal.
   $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau \nu_\tau)$ in a Flipped 2HDM with anomalously enhanced charged Higgs coupling to τ. [Phys. Rev. D 93, 115009]

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   (a) Explaining the observed deviation in R($D^{(*)}$) in an anomalous 2HDM.
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   (b) New tensor interaction as the source of the observed CP asymmetry in $\tau \to K_S \pi \nu_\tau$.
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12 to 16 December 2016, University of Delhi(oral presentation).

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Dedication

To mom, dad and uncle Tenpa Dorjee.

Acknowledgements

I want to thank Professor Nita Sinha for innumerable discussions and pointing out many logical blind spots and conceptual mistakes while I was doing my works toward this doctoral dissertation. Also I would like to thank Professor Nita Sinha for encouraging me to attend the CKM2016 conference at TIFR as well as the XXII DAE-BRNS High Energy Physics Symposium 2016 at Delhi University. I am also thankful to ma'm for writing me lots of recommendation letters for the many PDFs that I have applied. I would like to thank Rahul Srivastava for many discussions and comments about my work during his stay at IMSc as PDF. I would like to thank Professor D. Indumathi for many encouraging words and lending me Rs. 20,000 during the demonetization when I have to attend the CKM2016 and XXII DAE-BRNS High Energy Physics Symposium 2016. I am also thankful to ma'm for writing me lots of recommendation letters for the many PDFs that I have applied. I would like to thank Professor Shrihari Gopalakrishna for very helpful comments about the non-renormalizability of my model, in the now published work and also part of this thesis, regarding $B \to D^{(*)}\tau\nu_\tau$. I am also thankful to sir for writing me lots of recommendation letters for the many PDFs that I have applied. I would like to thank Professor G. Rajasekaran for writing me lots of recommendation letters for the many PDFs that I have applied. I would like to thank Professor M.V.N. Murthy for many encouraging words when he was in my doctoral committee before he retired. I would like to thank Professor Ghanashyam Date for many encouraging words when he was the chairman of my doctoral committee. In general, I would like to thank IMSc for providing me with many opportunities, good facilities and most important of all, for providing me with a great intellectual environment. I would also like to extend my gratitude to TCV, my alma-mater, which has given me so many great teachers who saw the seed in me to embark upon this great intellectual quest.

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SYNOPSIS OF Ph. D. THESIS

1. Name of the Student: DHARGYAL
2. Constituent Institution: Institute of Mathematical Sciences
3. Enrollment Number: PHYS10201104001
4. Title : PHENOMENOLOGICAL STUDIES OF THE OBSERVED ANOMALIES IN THE τ SECTOR.
5. Board of Studies: Physical Sciences

SYNOPSIS

Flavor is the mechanism by which the fermions in the Standard Model (SM) are distinguished from each other. In SM, this mechanism is intimately related to Higgs mechanism, because different fermions in the SM have different strength of Yukawa coupling to the Higgs particle and hence different masses. Also in SM, the source of the CP violation is related to the Higgs mechanism via the Kobayashi-Maskawa (KM) mixing matrix. But in general we can say that one of the least understood parts of the SM is the flavor sector of the SM, e.g why is the top quark mass much heavier than that of the say u quark? Why does the Higgs have such varied strength of couplings to the fermions while all the gauge bosons seem (to a very good precision) to have same coupling strength to all three generations? Flavor measurements provide very sensitive ways to probe SM and possible effects of New Physics (NP) if its effects can be observed at the LHCb or Belle -II or other current and future colliders, and so it can then put strong constraint on NP models. Currently there are several measurements showing deviations from SM predictions in the range of (2–4)σ level. In the following paragraphs, we will give a brief chapter wise summary of this thesis.

In chapter I, we will give a brief introduction of the key concepts underpinning the SM of particle physics, like Yang-Mills principle of local gauge invariance, spontaneous symmetry breaking and Higgs mechanism.

In chapter II, we will give a brief review of the flavor physics, especially related to the experimentally observed anomalies by Babar, Belle and LHCb in recent years. Currently there are several measurements showing deviations from SM predictions in the range of (2–4)σ level. In the following we will enumerate the most significant ones briefly.

$R_K = \frac{Br(\bar{B}\to\bar{K}\mu^+\mu^-)}{Br(\bar{B}\to\bar{K}e^+e^-)}$: The LHCb collaboration announced a 2.6σ deviation in the measurement of the ratio of the branching fraction of $\bar{B}$ to $\bar{K}$ and dimuons to that of $\bar{B}$ to $\bar{K}$ and dielectrons [50].

Muon $(g-2)$: The muon $(g-2)$ measurement has been reported to deviate from the SM expectation by more than 3σ though there still are quite large uncertainties in the SM predictions of this process.

$R(D^{(*)}) = \frac{Br(B\to D^{(*)}\tau\nu_\tau)}{Br(B\to D^{(*)}l\nu_l)}$: BaBar [20], Belle [21] and LHCb [22] measurements of $R(D^{(*)})$ decays with respect to SM predictions for these decays shows about 4σ (if we take the three deviations together). One of the main topics of this thesis is concerned with explaining the $R(D^{(*)})$ anomaly. Also there is the Babar collaboration's reported anomaly [4] in the measurement of CP asymmetry in $\tau \to K_S\pi\nu_\tau$ in the time integrated decay rate.

Chapter III focuses on the "Possible Hints of Lepton Flavor Universality Violation in R(D$^{(*)}$)". In the following paragraphs we will give a summary of the key results and conclusions [27]. BaBar [20], Belle [21] and LHCb [22] have reported an excess in the measurements of $R(D^*)$, $R(D)$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ than expected from SM, a possible signature of lepton flavor universality violating NP. In this work we have analyzed the implications for these decay modes in a Flipped/Lepton-Specific two Higgs doublet model (2HDM) with enhanced Yukawa coupling of $H^\pm$ to τ lepton with an enhancement factor of η. In general 2HDM the Yukawa Lagrangian can be writen as [29]:

$$\mathcal{L}^{2HDM}_{Yukawa} = -\sum_{f=u,d,l} \frac{m_f}{\mathcal{V}_0}(\xi^f_h \bar{f} f h + \xi^f_H \bar{f} f H - i\xi^f_A \bar{f}\gamma_5 f A)$$

$$-[\frac{\sqrt{2}V_{ud}}{\mathcal{V}_0}\bar{u}(m_u\xi^u_A P_L + m_d\xi^d_A P_R)dH^+ + \frac{\sqrt{2}\xi^l_A m_\tau}{\mathcal{V}_0}\bar{\nu_L} l_R H^+ + h.c.] \quad (1)$$

where $\xi^{q,l}$s depend on the type of 2HDM being used and $\mathcal{V}_0 = 246$ GeV is the vacuum expectation value of the Higgs field. But since we require a constructive interference of SM and charged Higgs contribution to fit all three of $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_\tau)$, only Lepton specific and Flipped 2HDM can achieve it. But contributions to the interference from just Lepton specific or Flipped 2HDM turn out to be too small to fit the three data simultaneously since the $\tan\beta$ dependence between m_b and m_τ cancels out in these models although they give constructive interference with SM unlike Type-I and Type-II 2HDM. Some additional factor has to be introduced into these models so that it can fit the three observables, R(D$^{(*)}$) and R($B \to \tau\nu_\tau$) simultaneously. The simplest and most straight forward way to achieve this enhancement is if we require that the τ lepton is screened from interacting with the full strength to the Higgs VEV v_2 and to all the neutral excitations from Higgs vacuum like the scalars h, H^0 and the pseudo-scalar A^0 in Flipped 2HDM. Since the Yukawa sector of the τ lepton breaks the $SU_L(2)$ symmetry the theory is not renormalizable in itself, so it has to be embedded inside a larger model in which Flipped 2HDM comes out as a 400 GeV–few TeV scale effective theory. The details of how such an anomalous interaction can be embedded inside a complete model is outside the scope of the present work and will be assigned to a future work. Here we focus on the phenomenological consequences of such a screening on observed B decay anomalies. Now since τ lepton is screened from seeing the full depth of v_2, its Yukawa coupling must increase so that its mass which is proportional to $Y_{Yukawa} \times \frac{v_2}{\eta}$ now is the same as the observed mass m_τ, which will effectively enhance the Yukawa coupling of τ lepton to the charged Higgs by a factor of η while Yukawa coupling of τ to h, H^0 and A^0 remains same as in usual Flipped 2HDM as the η factors in the neutral scalar Yukawa interaction cancel out with that of the η factor coming from reduced Higgs vacuum $\mathcal{V}_0$. Then the ξ_A^f factors in the charged Higgs interactions in this anomalous Flipped 2HDM with the screening factor η is given by:

$$\xi_A^u = \cot\beta,\ \xi_A^d = \tan\beta \text{ and } \xi_A^{e,\mu} = -\cot\beta \text{ and } \xi_A^\tau = -\eta\cot\beta. \tag{2}$$

Finally the charged Higgs Lagrangian contributing to $b \to cl\nu$ in this model can be written in the most general form using Yukawa couplings derived from ξ_A^fs and is given as:

$$\mathcal{L}_{H^\pm}^{Yukawa} = -[V_{cb}\bar{c}(g_s + g_p\gamma_5)bH^+ - \bar{\nu_l}(f_s^l + f_p^l\gamma_5)lH^+ + H.C] \quad (3)$$

where

$$\begin{aligned} g_s &= \frac{(m_b \tan\beta + m_c \cot\beta)}{\sqrt{2}\mathcal{V}_0}, g_p = \frac{(m_b \tan\beta - m_c \cot\beta)}{\sqrt{2}\mathcal{V}_0} , \\ f_s^{e,\mu} &= f_p^{e,\mu} = \frac{m_{e,\mu}\cot\beta}{\sqrt{2}\mathcal{V}_0} , \text{ and } f_s^\tau = f_p^\tau = \frac{m_\tau \cot\beta}{\sqrt{2}\mathcal{V}_0}\eta . \end{aligned} \quad (4)$$

Note the relative negative sign between the hadronic current and leptonic current which will ensure constructive interference between SM and charged Higgs contributions. In this model, SM and charged Higgs interfere constructively, unlike Type-I and Type-II 2HDM models where they interfere destructively. In this modified Flipped 2HDM, the charged Higgs coupling to τ is enhanced by a factor of η than that of a simple Flipped 2HDM. Although the η is an independent parameter but if we require $\eta = \tan^2\beta$, then the Yukawa interactions in the hadronic sector of our model has the same form as in the Type-II 2HDM but interaction in the leptonic sector of our model has the same form as in Flipped 2HDM except the τ lepton coupling to the charged Higgs, which has same form as in Type-II 2HDM but with opposite sign. In Table 1 we have shown three different values of the parameters $\tan\beta$ and $M_\pm$ that fits at same accuracy. From Table 1, we see that

S.no	$\tan\beta$	$M_\pm$ GeV	$R(D)_{Th}$	$R(D^*)_{Th}$	$Br_{Th}(B \to \tau\nu)$
1	39.98	400	0.348	0.255	1.29×10^{-4}
2	69.97	700	0.348	0.255	1.29×10^{-4}
3	99.95	1000	0.348	0.255	1.29×10^{-4}

Table 1: $\chi^2_{min} = 10.95$ and we have restricted the $\tan\beta$ in the range of $100 > \tan\beta > 1$.

in the range 1 TeV $\geq M_{H^{\pm}} \geq$ 400 GeV and $100 > \tan\beta > 1$, we have:

$$R(D)_{Th} = 0.348 \pm 0.16 \quad (5)$$

$$R(D^*)_{Th} = 0.255 \pm 0.07 \quad (6)$$

and

$$Br_{Th}(B \to \tau\nu) = (1.29 \pm 0.89) \times 10^{-4}, \quad (7)$$

compared to the combined [Babar,Belle,LHCb] [23] experimental values:

$$R(D)_{EXP} = 0.388 \pm 0.047 \quad (8)$$

$$R(D^*)_{EXP} = 0.321 \pm 0.021 \quad (9)$$

and

$$Br_{EXP}(B \to \tau\nu) = (1.14 \pm 0.27) \times 10^{-4}. \quad (10)$$

By adding theoretical and experimental errors in quadrature from Eqs(5,6,7) and Eqs(8,9,10) respectively, we conclude that our phenomenological model can give results in agreement within 1σ deviation for the combination of $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ compared to about 4σ deviation from SM for the latest combined [Babar,Belle,LHCb] experimental data for these observables. The same results can be achieved if b quark replaces the τ lepton in a Lepton Specific 2HDM. In that case Yukawa coupling of the leptonic sector will be same as Type-II 2HDM and Yukawa coupling of the quark sector will be same as Lepton Specific 2HDM [29] except the b quark which will have the effective Yukawa coupling same as in Type-II with opposite sign. In that case charged Higgs mass may be allowed to be lower than 400 GeV. We also observed that from the form of the Yukawa couplings it is expected that if we require $\eta = -1$ for the b quark or τ lepton in the 2HDM-II, then also it will fit the three observables at about same accuracy as above model of Flipped/Lepton-Specific 2HDM. This is interesting in a sense that it will be like 2HDM-II but with a wrong sign in either b quark or τ lepton Yukawa coupling.

In Chapter IV, we will give a review of the theoretical analysis of CP violation in hadronic τ decays as well the experimental status of the some of the key CP observables in hadronic τ decays.

Chapter V contains another main topic of this thesis "Observed Sign Anomaly in $A_{cp}(\tau \to K_S \pi \nu_\tau)$". In the following paragraphs we will give a summary of the key results and conclusions of our analysis of the contributions coming from tensorial current to this observable [8][19]. Babar collaboration has reported an intriguing opposite sign in the time integrated decay rate asymmetry in $\tau \to K_S \pi \nu_\tau$ than that of SM prediction from the known K^0 - $\bar{K}^0$ mixing. Babar's results [4] deviate from the SM prediction by about 2.7σ. In [8], we have shown that the CP violation coming from the $K - \bar{K}$ mixing and the direct CP violation in $A_{cp}(\tau \to K_S \pi \nu_\tau)$ can be separated as

$$A_{cp}(\tau \to K_S \pi \nu_\tau) = \frac{A^K_{cp} + A^\tau_{cp}}{1 + A^K_{cp} A^\tau_{cp}} \ , \tag{11}$$

where A^K_{cp} is the CP violation arising from the $K - \bar{K}$ mixing and A^τ_{cp} is the direct CP violation arising from a NP particle mediating CP violation at the lepton and/or quark vertices and τ_τ is the τ life time. Since both A^K_{cp} and A^τ_{cp} are expected to be small, we can safely ignore terms involving the product of the two in the denominator. In [8], using the Breit Wigner forms of the various resonances contributing to the vector and scalar form factors, the $s(K_S\pi)$ (hadronic invariant mass) dependent strong phases were determined. Taking the strong phase of the tensor form factor to be vanishing and its magnitude to be constant, the observables $A_{cp}(\tau \to K_S \pi \nu_\tau)$ and branching fraction were used to determine weak phase and the magnitude of the tensor coupling. Also in this work we have given the full angular distribution of the decay rate and CP asymmetry in presence of new tensor current of the $\tau \to K_S \pi \nu_\tau$.

Publications in Refereed Journals:

1. Published:

   (a) H. Zeen Devi, L. Dhargyal and Nita Sinha, "Can the observed CP asymmetry in $A_{CP}(\tau \to K_S \pi \nu)$ be due to nonstandard tensor interactions?" [Phys.Rev.D90(2014)no.1,013016]

   (b) Lobsang Dhargyal, "$R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ in a Flipped 2HDM with anomalously enhanced charged Higgs coupling to τ." [Phys. Rev. D 93, 115009 (Published 7 June 2016)]

2. Communicated:(not included in the thesis)

   (a) Lobsang Dhargyal, "Full angular spectrum analysis of tensor current contribution to $A_{cp}(\tau \to K_S \pi \nu_\tau)$." [arXiv:1605.00629 (2016)]

3. Conference Presentations:

   (a) Explaining the observed deviation in R($D^{(*)}$) in an anomalous 2HDM.
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4.	**D. Indumathi**	Member		
5.	**Srihari Gopalakrishna**	Member		
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3. Following papers are not part of the thesis :

   (a) Lobsang Dhargyal.
   "Search for more sensitive observables to charged scalars in $B \to D^{(*)}\tau\nu_\tau$." [arXiv:1703.10735v2 [hep-ph] 10 April 2017]

   (b) Lobsang Dhargyal.
   "A simple extension of SM that can explain the $(g-2)_\mu$ anomaly, small

neutrino mass and dark-matter." [arXiv:1705.09610 [hep-ph] 24 May 2017]

4. <u>Conference Presentations:</u>

   (a) Explaining the observed deviation in R($D^{(*)}$) in an anomalous 2HDM.
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Chapter 1

Introduction

Classical view:

"I shall be telling this with a sigh
Somewhere ages and ages hence:
Two roads diverged in a wood, and I—
I took the one less travelled by,
And that has made all the difference."
by Robert Frost.

Quantum view:

"I shall be telling this with a sigh
Somewhere ages and ages hence:
Two roads diverged in a wood, and I—
I took both of them simultaneously,
And that has made all the difference."
by Lobsang Dhargyal.

Towards a synthetic view?

"Whatever degenerations there are in the world,
The root of all these is ignorance;
You taught that it is dependent origination,
The seeing of which will undo this ignorance.

....

Intrinsic nature, non-constructed and non-contingent,
Dependent origination, contingent and constructed—
How can these two converge
Upon a single basis without contradiction?
Therefore whatever originates dependently,
Though primordially free of intrinsic existence,
Appears as if it does (posses intrinsic existence);
So you taught all this to be illusion-like."

by Je Tsongkhapa.

Since the time of Galileo and Issac Newton, the key people who more or less have the claim of initiating the scientific revolution, our view and understanding of natural world and its processes have seen great progress. Our theories regarding the natural world have come all the way, from a static and independent space and time background stage upon which the great cosmic dance by constituent particles and fields are performed, to the special relativity, where space and time lose their independence from each other as well as from relative motions of observers and it only makes sense to promote space and time on equal footing as space-time, and to the general relativity where space-time itself becomes dynamical and participates in the laws governing the motions of particles and fields and can be affected and changed by matter and energy contained in it. Then came the most drastic conceptual shift from the classical world view, the Quantum revolution, beginning with Max Planck, Einstein and Bohr's quantum theory to Born, Heisenberg, Schroedinger and Dirac formulations of Quantum mechanics, which incorporates the experimental findings that particles at atomic and smaller scales start showing both wave like behavior as well as particle like behavior. Then came the Quantum Field Theory (QFT), which is a synthetic product of merger of special relativity and quantum mechanical principles. In QFT, the quantum dynamical laws, the commutator relations, are implemented in the expansion coefficients of the irreducible representation bases of Poincaré group, special relativity, as a background field which carries the space-time information of the particle-field; the particle-like behavior, the commutator relations, are carried by the generalized Fourier coefficients and the Hilbert bases wave function carries the wave-like behavior of the particle-fields. In essence, QFT is a mathematical construct where the two contradictory natures of particle-like behavior and wave-like behavior of particle-fields are put together covariantly. The relativity, quantum mechanics and QFT are the foundation stone upon which the modern formulation of Standard Model (SM) of particle physics is based.

The Standard Model of Particle Physics

The SM of particle physics encompasses our accumulated knowledge about the laws governing the natural world, except the very large scale structure which is governed by gravity. It is a renormalizable QFT model based on the Yang-Mills principle of non-abelian local gauge invariance [49], spontaneous symmetry breaking and Higgs mechanism. The SM of particle physics is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, where the $SU(3)_C$ is the local gauge group governing the strong interaction and $SU(2)_L \times U(1)_Y$ is the local gauge group governing the electro-weak interaction. The gauge interactions are mediated by the gauge bosons, where the gauge boson mediating the strong interactions are the eight gluons and the gauge bosons mediating the electro-weak interaction are the photon, which mediate electromagnetic force, and the weak gauge bosons, $W^{\pm}$ and Z, which mediate the weak interaction.

Yang-Mills Principle of Local Gauge Invariance

The Yang-Mills principle of non-abelian local gauge invariance is the generalization of the basic idea that if we make the parameter of the global gauge transformation dependent on space-time, then the terms which contain derivatives and terms which do not contain derivative in the Lagrangian will behave differently. And to make the Lagrangian invariant under local gauge transformations, new degrees of freedom are introduced, the gauge particles, which will transform under the local gauge such that the free field and its interaction with the new gauge fields will be invariant as a whole, even though each of them individually are not invariant under the local gauge transformation. The particle content of the SM can be classified according to their transformations under the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ as:

$$L_L(1,2,-1) = \begin{pmatrix} \nu_L \\ l_L \end{pmatrix}, \ l_R(1,1,-2) \ , \tag{1.1}$$

where $L_L(1,2,-1)$ and $l_R(1,1,-2)$ are the left and right handed $SU(2)_L$ doublet and singlet leptons respectively and

$$Q_L(3,2,+1/3) = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \ u_R(3,1,+4/3), \ d_R(3,1,-2/3) \ , \tag{1.2}$$

where $Q_L(3,2,+1/3)$ are the left handed $SU(2)_L$ quark doublets and $u_R(3,1,+4/3)$ and $d_R(3,1,-2/3)$ $SU(2)_L$ singlet quarks. There are three families for each of the leptons and quarks above with different masses; this structure is called flavor structure in general, and it is clear that the flavor structure is very intimately related to the mass generation mechanism, as the only thing that differentiates flavor of leptons and quarks are their different masses. So the particle content of SM are 8 gluons mediating the strong interaction, one photon mediating the electromagnetic interaction, three weak gauge bosons mediating the weak interaction, three up type quarks (u, c, t), three down type quarks (d, s, b), three charged leptons (e, μ, τ) and three left handed neutrinos $(\nu_e, \nu_\mu, \nu_\tau)_L$. But it is borne out of experiment that neutrinos have mass even though their masses are much smaller than that of the charged fermions, so if we include the neutrino mass into account then the simplest extension of SM is to introduce a right handed neutrino and give them a Dirac masses via Higgs mechanism similar to the charged particles νMSM (neutrino-minimum-SM), but then we do not understand why neutrino masses should be so small compared to the charged fermion masses.

Spontaneous Symmetry Breaking And Higgs Mechanism

The gauge invariance condition forbids the gauge bosons to have mass terms explicitly in the Lagrangian, so the idea of spontaneous symmetry breaking (SSB) was adapted by Y. Nambu into the particle physics (HEP) from the condensed matter physics (LEP). The basic idea of SSB is that we introduce a new scalar particle into the Lagrangian which develops a non zero vacuum expectation value (VEV), which breaks the gauge symmetry in such a way that the Lagrangian as a whole is still symmetric under the gauge group, but the vacuum of the scalar does not follows the symmetry of the Lagrangian. The Higgs mechanism is one particular model of SSB which turns out to be a correct (nominally) model, after the discovery of SM Higgs like scalar by LHC in 2012, to give masses to the electroweak gauge bosons as well as to the fermions. The part of the Lagrangian which expresses the Higgs particle's self interaction and its interaction with the electroweak gauge bosons, which provides a mechanism to give masses to the gauge bosons without breaking the gauge symmetry of the Lagrangian—i.e., conserved charges due to gauge symmetries are still conserved—is given as:

$$L_\Phi = (D_\mu \Phi)^\dagger D^\mu \Phi - \lambda(|\Phi|^2 - \mathcal{V}_0^2/2)^2 + \lambda \mathcal{V}_0^4/4 \ , \tag{1.3}$$

where, in the unitary gauge, we have,

$$\Phi = \frac{1}{\sqrt{2}} \left[\begin{array}{c} 0 \\ \mathcal{V}_0 + H \end{array} \right] \tag{1.4}$$

with $D^\mu \Phi = (\partial^\mu + i(g)\sigma_i W_i^\mu + ig'/2Y_\Phi B^\mu)\Phi$ which gives

$$(D_\mu \Phi)^\dagger D^\mu \Phi \to 1/2\partial_\mu H \partial^\mu H + [g^2/4W_\mu^\dagger W^\mu + \frac{g^2}{8\cos^2\theta_W} Z_\mu Z^\mu](\mathcal{V}_0 + H)^2 \ . \tag{1.5}$$

This give the gauge boson masses as

$$M_W = M_Z \cos\theta_W = \frac{1}{2}\mathcal{V}_0 g \ . \tag{1.6}$$

This is the well known Higgs mechanism to generate mass for the gauge bosons via spontaneous symmetry breaking.

Fermion Masses and Flavor Structure of SM

The Yukawa interaction of the Higgs doublet with the fermions in the SM is given as

$$L_Y = -\sum_{i,j=1}^{3} \left[\bar{Q}_{Li}(y_{ij}^d \Phi d_{jR} + y_{ij}^u \widetilde{\Phi} u_{iR}) + \bar{L}_{Li} y_{ij}^l \Phi l_R \right] + h.c. \ , \tag{1.7}$$

which when substituted for the Higgs doublet from Eq. (1.4) gives

$$L_Y = -\left[\bar{u}_L M_u' u_R + \bar{d}_L M_d' d_R + \bar{l}_L M_l' l_R + h.c \right] \left(1 + \frac{H}{\mathcal{V}_0} \right) \ , \tag{1.8}$$

where M_u', M_d' and M_l' are the mass matrices of the up type quarks, down type quarks and charged leptons respectively. Now when we diagonalize the M's by multiplying from left by U_L and from right by U_R, we get diagonal mass matrices given as

$$L_Y = -\left[\bar{u} M_u u + \bar{d} M_d d + \bar{l} M_l l \right] \left(1 + \frac{H}{\mathcal{V}_0} \right) \ , \tag{1.9}$$

with M's being diagonal and real. These unitary transformations will not affect the neutral current interaction, but it will affect the charged current interactions by introducing the Cabibbo-Kobayashi-Maskawa (CKM) matrix element multiplying each of the respective charged currents given as

$$L_{CC} = -\frac{g}{2\sqrt{2}} \left\{ \left[\sum_{ij} \bar{u}_i \gamma^{\mu}(1-\gamma_5) V_{ij} d_j + \sum_{ij} \delta_{ij} \bar{\nu}_i \gamma^{\mu}(1-\gamma_5) l_j \right] W^{\dagger}_{\mu} + h.c \right\} \ , \quad (1.10)$$

where the V_{ij} are the CKM matrix elements given by

$$V = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \quad (1.11)$$

which satisfies the unitary conditions of $VV^{\dagger} = V^{\dagger}V = 1$; its elements are to be determined experimentally. For 3×3 unitary matrices, there will be three rotation angles and one irreducible phase, this phase is the only source of violating the combined symmetries of Charge-Conjugation and Parity (CP) in the SM.

Chapter 2

Brief Review of Flavor Physics

Introduction

Flavor is the mechanism by which different fermions in the SM distinguish one from the others. In SM, this mechanism is intimately related to Higgs mechanism, because different fermions in the SM have different strength of Yukawa coupling to the Higgs particle and hence different masses accordingly. If all the fermions in the SM are massless then SM has $U(3)^3$ flavor symmetry in the quark sector and $U(3)^2$ flavor symmetry in the lepton sector, where we have $U(3)_Q \times U(3)_u \times U(3)_d$ in quark sector and $U(3)_L \times U(3)_l$ in lepton sector. This flavor symmetry in the SM is broken by Yukawa interaction of the SM fermions to the Higgs field. Also in SM, the source of the CP violation is related to the Higgs mechanism via the KM (Kobayashi-Maskawa) mixing matrix. In general we can say that one of the least understood part of the SM is the flavor sector of the SM, e.g., why is the top quark much heavier than say the u quark? Why does the Higgs have such a varied strength of coupling to the fermions while all the gauge bosons seem (to a very good precision) to have the same coupling strength to all three generations? etc. In any case, flavor measurements provide very sensitive ways to probe SM and possible effects of NP if there are reachable by future LHCb and Belle-II sensitivity, and so it can put strong constraint on NP models. However, currently there are several measurements showing deviations from SM predictions in the range of (2–4)σ level. In the following sections, we will give a brief review of those measurements, some of which could be just statistical fluctuations, but some may be hints of NP.

Flavor And CP Violation In The SM

In the SM the one irreducible phase in the CKM matrix in the quark sector is the only source of CP violation besides the CP violation in the strong interaction which turns out to be a very small effect (the strong CP problem). In the Wolfenstein parametrization, we have the CKM matrix elements given as:

$$V_{CKM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_b \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} = \begin{pmatrix} 1-\lambda^2/2 & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1-\lambda^2/2 & A\lambda^2 \\ A\lambda^3(1-\rho-i\eta) & -A\lambda^2 & 1 \end{pmatrix} + O(\lambda^4) \tag{2.1}$$

where the expansion parameter λ is extracted from experiments in flavor observables and given to be $\lambda \approx 0.23$.

Since CKM matrix is unitary we have

$$\sum_k V_{ik}V^*_{jk} = \sum_k V_{ki}V^*_{kj} = \delta_{ij} \tag{2.2}$$

and the six vanishing relations from the above can be represented by triangles in a complex plane with all the triangles having the same area. The most famous of these triangles is the one extracted from the relation

$$V_{ud}V^*_{ub} + V_{cd}V^*_{cb} + V_{td}V^*_{tb} = 0. \tag{2.3}$$

The Figure 2.1 shows the triangle from the above relation where the angles of the triangle in the same figure are given as

$$\alpha = arg\left(-\frac{V_{td}V^*_{tb}}{V_{ud}V^*_{ub}}\right),\ \beta = arg\left(-\frac{V_{cd}V^*_{cb}}{V_{td}V^*_{tb}}\right),\ \gamma = arg\left(-\frac{V_{ud}V^*_{ub}}{V_{cd}V^*_{cb}}\right). \tag{2.4}$$

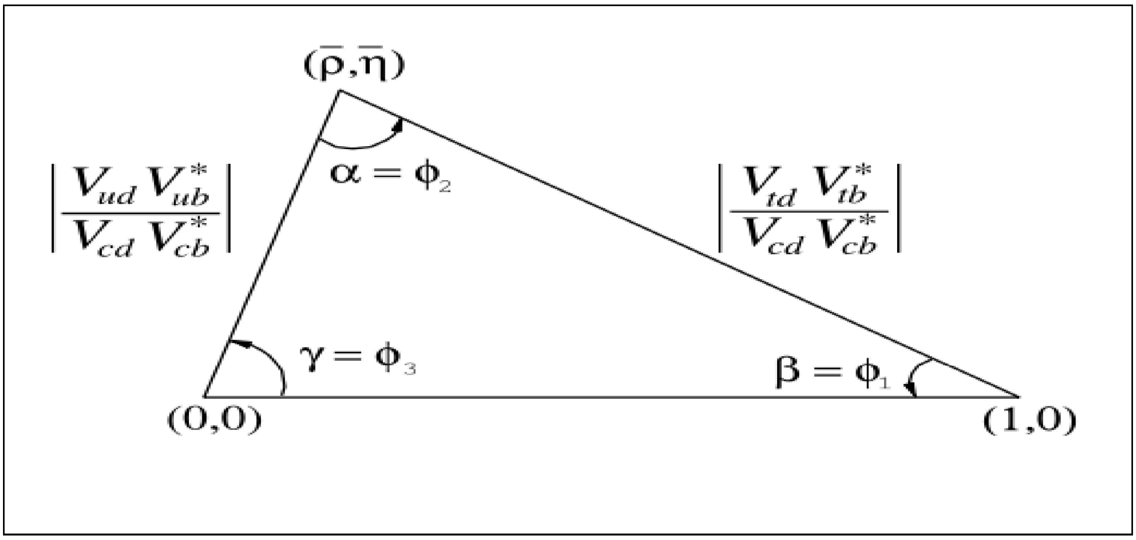

Figure 2.1: Unitary triangle from Eq. (2.3) [66].

The magnitudes of the sides of CKM triangles are related to the sizes of the unitary triangle. The magnitude and phase of CKM elements are mainly extracted from semi-leptonic and leptonic K and B decays and $B_{d,s}$ mixing. The measurement of the angles of the unitary triangle and therefore non-zero area of the triangle involves measurements of CP violation, as only in CP violating observables the phase in CKM shows its effect. In Figure 2.2 we have shown the SM CKM fit of the parameters $\bar{\eta}$ and $\bar{\rho}$ (where the $\bar{\eta}$ and $\bar{\rho}$ are η and ρ with corrections made on the parameters to preserve the unitarity of the CKM triangle) [66], and the best experimental constraints on the magnitudes and angles of the triangle.

Observed experimental anomalies in flavor physics at $(2–4)\sigma$ level

R_K : The LHCb collaboration announced a 2.6σ deviation in the measurement of the ratio of the branching fraction of B decaying into di-muons to that over di-electrons [50],

$$R_K = \frac{Br(B \to K\mu^+\mu^-)}{Br(B \to Ke^+e^-)} \ , \tag{2.5}$$

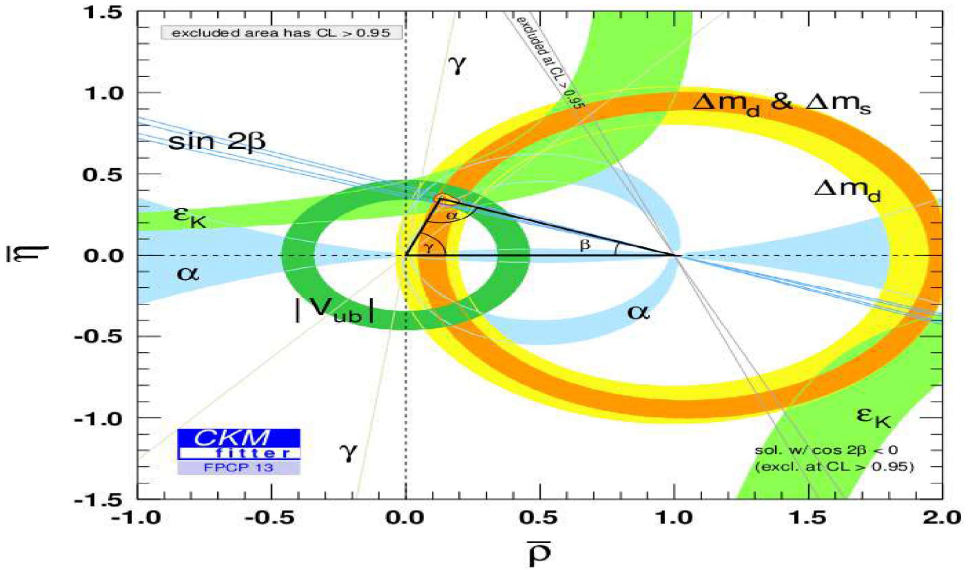

Figure 2.2: The SM CKM fit and experimental constraints on them [66].

and obtained experimental value of above ratio as

$$R_K = 0.745 \pm^{0.090}_{0.074} \pm 0.036\ , \tag{2.6}$$

which shows deviation of 2.6σ from the SM prediction of the $R_K = 1.0003 \pm 0.0001$ including the α_S and subleading $1/m_b$ corrections. Also the latest LHCb measurements of $\frac{Br(B \to K^* \mu^+ \mu^-)}{Br(B \to K^* e^+ e^-)}$ give [64]

$$R_{K^*} = 0.66 \pm 0.111(Stat) \pm 0.03(Sys) \text{ for } 0.045 < q^2 < 1.1 \text{ GeV}^2\ , \tag{2.7}$$

and

$$R_{K^*} = 0.69 \pm 0.111(Stat) \pm 0.05(Sys) \text{ for } 1.1 < q^2 < 6 \text{ GeV}^2\ , \tag{2.8}$$

which shows deviation of 2.4σ and 2.6σ respectively from the SM prediction. There could be many possible connections between this anomaly and the P_5' anomaly observed in the $B \to K^* \mu^+ \mu^-$ in angular distribution [51] measured by LHCb [52] and Belle [53]. Another anomaly that could be related to above two is the reported measurement of the branching fraction ratio of $B_S \to \phi\mu^+\mu^-$, which is about 3σ lower than expected from theoretical calculations [54]. If NP are to contribute to these modes then it is very likely to impact $B \to \mu^+\mu^-$ too, and it could place very strong constraint on the NP models. Another 3–4σ anomaly (deviation from SM expectation) is in the $D\emptyset$ measurement of the like-sign dimuon charge asymmetry in semileptonic decays of b hadrons, $(N_{\mu^+\mu^+} - N_{\mu^-\mu^-})/(N_{\mu^+\mu^+} + N_{\mu^-\mu^-})$[55]. But recent measurements of BaBar, Belle, $D\emptyset$ and LHCb seem to be consistent with SM prediction.

There is also the reported excess in $Br(h \to \tau\mu) = (0.84^{+0.39}_{-0.37})\%$ by CMS [56] and ATLAS [57] of about 2.4σ when the excess from CMS and ATLAS are combined. This is interesting in a sense that in SM this decay is not allowed by conservation of lepton flavor in the charged leptonic sector but many extensions of SM allow lepton flavor violation (LFV) dynamics in the charged lepton sector.

Muon $(g-2)/2$: The muon $(g-2)/2$ measurements have been reported to deviate from the SM expectation by more than 3σ, however there is still uncertainty in the SM calculations as there are complicated strong interaction effects involved in the calculations, but it could also be due to effects from NP. Once reliable lattice QCD calculations of the hadronic light by light scattering and vacuum polarization contributions are known then only we can know whether it is due to NP or due to some complicated strong interaction dynamics. Still supersymmetric models with light sleptons can account for the deviation [63].

$R^W_{\tau l}$: The long standing LEP2 measurement of [58]

$$\frac{2Br(W \to \tau\bar{\nu}_\tau)}{Br(W \to e\bar{\nu}_e) + Br(W \to \mu\bar{\nu}_\mu)} = 1.055 \pm 0.023, \tag{2.9}$$

indicates an anomaly of about 2.4σ from expectation of the lepton universality of the SM gauge couplings $R^W_{\tau l}|_{SM} = 0.999$ [60] with negligible error. This measurement received not much attention from the NP model builders as a simple modification of W vertex to accommodate this anomaly will drastically effect the per mil level agreement of lepton universality in tau, mesons and Z decays.

There is also the long standing tension between the inclusive and exclusive measurements of $|V_{cb}|$ and $|V_{ub}|$. The $|V_{cb}|$ and $|V_{ub}|$ together with the γ determine the apex of the unitary triangle from tree level processes, which plays a crucial role in the improvement of the sensitivity to NP in B mixing and NP contribution to CP violation in the loop processes.

$R(D^{(*)})$: The radiative penguin decays and tauonic B decays provide excellent probes in search of NP. Recent measurements in these sectors of flavor physics by BaBar, Belle and LHCb showed impressive experimental precision, providing strong constraints on NP models. Most of the measurements agree with SM but quite unexpected and large deviations from SM have been reported, surprisingly in decays that occur at tree level in SM. The reported tension in the measurements by BaBar, Belle and LHCb measurements of $Br(B \to \tau\nu_\tau)$ and $R(D^{(*)})$ decays with SM predictions for these decays shows about 4σ if we take the three deviations together; this is the most significant discrepancy from the SM in collider experiments so far (aside from the non zero neutrino masses). The Type-II 2HDM has been ruled out by BaBar[61] and Belle[62] at 99.8% confidence level. In Fig. 3.1 we have shown the BaBar's plot of $R(D)$ and $R(D^*)$ in the context of the Type-II 2HDM theoretical predictions as a function of $\frac{\tan\beta}{m_{H^+}}$. The measured values of the $R(D)$ and $R(D^*)$ match the predictions of the Type-II 2HDM for $\frac{\tan\beta}{m_{H^+}} = 0.44 \pm 0.02$ GeV^{-1} and $\frac{\tan\beta}{m_{H^+}} = 0.75 \pm 0.04$ GeV^{-1}, respectively. However the combination of $R(D)$ and $R(D^*)$ data excludes the contribution from a charged Higgs of Type-II 2HDM at 99.8% confidence level for any value of $\tan\beta/m_{H^+}$ as is clearly shown in Fig. 3.1.

Chapter 3

Possible Hints of Lepton Flavor Universality Violation in R(D$^{(*)}$)

3.1 Introduction

Decays of B mesons into final states containing τ lepton are sensitive to new charged current interactions that can induce lepton flavor universality violation. The decays $B \to D^{(*)}\tau\nu_\tau$ and $B \to \tau\nu_\tau$ are well suited for searching for effects of new physics (NP) in charged current interactions. In particular, the presence of third generation fermions in both initial and final states implies that these kind of decays are very sensitive to new physics (NP) related to the Higgs sector such as charged Higgs. The large data samples of the B-factories and recent advances in techniques of full-event reconstruction have led to the evidence of the decay $B \to \tau\nu_\tau$ and unambiguous observation of the decays $B \to D^{(*)}\tau\nu_\tau$. The multiple neutrinos produced in these exclusive decays make it very difficult to reconstruct the invariant mas of the B meson and use it for background rejection. Therefore, their study requires use of additional constraints related to the production of the B meson. Such constraints are easily available at B factories, which collide electrons and positrons at an average center of mass energy of $\sqrt{s} = 10.58$ GeV, corresponding to the mass of the $m_{\Upsilon(4S)}$ of the $\Upsilon(4S)$ resonance. As a result the B factories, BABAR and Belle, are the most suited for search for NP in these decays. During the next decade, the Belle-II experiment, which will have an integrated luminosity over 30 times greater than that of the combined data sets of Belle-I and BABAR, will provide very accurate measurements that can constrain possible new physics (NP) with great precision. This makes recent BABAR, Belle and LHCb [20, 21, 22] reporting of an excess in the measurements of $R(D^{(*)})$ and $Br(B \to \tau\nu_\tau)$ than expected from Standard Model (SM), a possible signature of lepton flavor universality violating new physics (NP), very interesting from the point of view of new physics (NP) searches.

3.2 The Technique of Full-Event Reconstruction

The technique of full event reconstruction has played and will play a very important role in the measurements of these decays in future B factories. A B factory is a high-luminosity e^+e^- collider with an average center-of-mass (CM) energy $\sqrt{s}$ equal to the $\Upsilon(4S)$ mass, i.e, $m_{\Upsilon(4S)} = 10.57$ GeV. In what follows we will take all kinematic quantities in the CM frame. The produced $\Upsilon(4S)$ decays promptly into two B pairs, so that the B energy equals to $\sqrt{s}/2$ to within half the collision energy spread, which is $\sigma_{\sqrt{s}/2} = 5$ MeV at current B factories. The momenta of the B mesons, which average about 330 MeV, are equal to within $\sigma_{\sqrt{s}/2}$ and opposite to within 2°. These collision event informations are essential for the technique of full event reconstruction.

These event characteristics are used to address the difficulties caused by the undetectable in rare B meson decays. This is done by reconstructing not only the signal B meson decays (labeled B_{sig}), but also the other B meson in the event, known as tag B (labeled B_{tag}). In such full event reconstruction, it is typically required that all charged particle tracks be assigned to one of the two B meson candidates. Also, the energy E_{extra} of the unassigned calorimeter clusters or the photon candidates is required to be low, typically around 1 GeV. This requirement reflects the fact that such extra energy arises not only from missing particles in the background events, but also from calorimeter noise, previous events, and scattered particles from the interaction of hadrons with the calorimeter in the signal events. By attempting to account for the origin of all the particles in the event, full event reconstruction reduces the rate of the combinatorial background, which arises from the random combinations of particles that happen to satisfy the selection criteria. Furthermore, if the tag B is fully and correctly reconstructed in a hadronic final sate, the kinematic constraints described above yield a measurement of the four momentum of the missing neutrinos, further aiding with the signal identification and enabling the calculation of the quantities in the signal B rest frame.

The disadvantage of full event reconstruction is the low efficiency for reconstructing the large number of particles produced in a typical tag B decay. Nevertheless, the large datasets of BABAR and Belle and increasing sophistication in the application of tag B reconstruction technique have made this technique an indispensable tool kit for the rare B decays and B decays with multiple neutrinos in the final state. Tag B reconstruction is done with three main techniques; hadronic tagging, semi-leptonic tagging, or inclusive, depending on the tag B final state and reconstruction method deployed. The different tagging methods are generally complementary, with each method having advantages and disadvantages, and relative importance depending on the signal B decay of interest among other things.

3.3 Experimental Measurements of $R(D^{(*)})$ and $Br(B \to \tau\nu_\tau)$

First Babar [20] and then Belle [21] have reported a possible hint of lepton flavor universality violating NP in $R(D^{(*)}) = \frac{Br(B\to D^{(*)}\tau\nu)}{Br(B\to D^{(*)}l\nu)}$. Recently LHCb [22] has also measured a deviation in $R(D^*)$ from SM. In Table 3.1 we have given the experimental measurements of $R(D^{(*)})$ [42, 43] and their averages [44] from BABAR, Belle and LHCb collaborations and the SM expectations for these observables [45]. The combined (Babar,Belle and

LHCb) results gives [23],

$$\begin{aligned} R(D)_{EXP} &= 0.388 \pm 0.047 \ , \\ R(D^*)_{EXP} &= 0.321 \pm 0.021 \ . \end{aligned} \tag{3.1}$$

Comparing these measurement with the SM predictions [24],

$$\begin{aligned} R(D)_{SM} &= 0.297 \pm 0.017 \ , \\ R(D^*)_{SM} &= 0.252 \pm 0.003 \ . \end{aligned} \tag{3.2}$$

There is a deviation of 1.8σ for the R(D) and 3.3σ for the $R(D^*)$ and the combination corresponds to a close to 3.8σ deviation from SM prediction.

Collaboration	R(D)	$R(D^*)$
BABAR	0.44±0.058±0.042	0.332±0.024±0.018
Belle	0.375±0.064±0.026	0.293±0.038±0.015
Belle		0.302±0.030±0.011
LHCb		0.336±0.027±0.030
Exp. average	0.397±0.040±0.028	0.316±0.016±0.010
SM expectation	0.30±0.010	0.252±0.005

Table 3.1: Experimental measurements of $R(D^{(*)})$ and their averages from BABAR, Belle and LHCb collaborations and the SM predictions [65].

It is further supported by measurement of Br($B \to \tau\nu$) by Babar [25] and Belle [26], with the PDG average:

$$Br_{EXP}(B \to \tau\nu) = (1.14 \pm 0.27) \times 10^{-4} \ , \tag{3.3}$$

which is 1.3σ above the SM prediction [28]. Now if we add the errors in $R(D^*)$ and $B(B \to \tau\nu)$ in quadrature then the three data combined deviates from SM by about 4σ.

Currently, the $R(D^*) = \frac{Br(B \to D^{(*)}\tau\nu)}{Br(B \to D^{(*)}l\nu)}$ ratios where l is e or μ constitute the most significant discrepancy from SM in collider experiments beside neutrino masses. The effect is at about 4σ level, which imply that if these anomalies persist in future measurements then NP contributing to these observables should be at a fairly low scale.

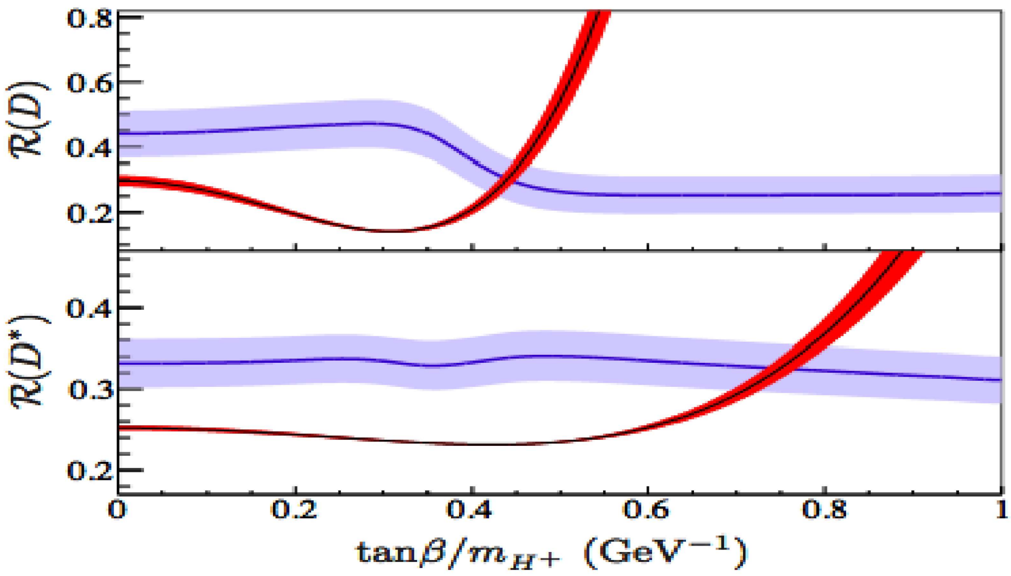

Figure 3.1: Plots of the BaBar's results (light blue band) compared to the predictions of charged Higgs effect from the 2HDM of Type-II (dark red band) as a function of $\frac{\tan\beta}{m_{H^+}}$ for the observables $R(D)$ and $R(D^*)$. SM corresponds to $\frac{\tan\beta}{m_{H^+}} = 0$ [20].

The effect of Physics beyond the SM in Tauonic B decays has been studied extensively, particular in the context of Type-II 2HDM but Babar [20] has ruled out Type-II 2HDM at 99.8 percent confidence level from the fits to the $R(D^{(*)})$ data for any values of $\frac{\tan\beta}{M_\pm}$. However, Belle's measurements shows some compatibility with Type-II 2HDM at about $\frac{\tan\beta}{M_\pm} = 0.5$ GeV^{-1} but if we include the $B \to \tau\nu_\tau$ data also then it turns out that the contribution from such a large $\frac{\tan\beta}{M_\pm}$ actually pushes the $B \to \tau\nu_\tau$ well over 5σ from the PDG world average. This is mainly due to the fact that since Type-I and

Type-II 2HDMs interfere destructively with SM as shown in Fig. 3.1, in which BaBar's results (light blue band) are plotted against the predictions of charged Higgs effect from the 2HDM of Type-II (dark red band) as a function of $\frac{\tan\beta}{m_{H^+}}$ for the observables $R(D)$ and $R(D^*)$. Fig. 3.1 clearly shows that at low $\frac{\tan\beta}{m_{H^+}}$, the interference term between SM and charged Higgs dominates and hence the the plot goes downwards but at high $\frac{\tan\beta}{m_{H^+}}$ the purely charged Higgs contribution get larger than the interference term and hence the plot goes upward, so it is clear that in these type of 2HDM only a large contribution from purely charged Higgs can have positive contribution, but a large contribution from purely charged Higgs part that gives a moderate push to the $R(D^{(*)})$ pushes the $B \to \tau\nu_\tau$ beyond 5σ from the world average very fast. So only a 2HDM with constructive interference with SM like Flipped or Lepton specific 2HDM may be able to fit all three of the B-physics data $R(D^{(*)})$ and $B \to \tau\nu_\tau$. But Flipped or Lepton specific 2HDM in itself does not fit because in these models the $\tan\beta$ dependence between m_b and m_τ in the interference between SM and charged Higgs cancels out and so there is no corresponding increase in contribution from interference part as $\tan\beta$ increases. To fit all three of the data we need to modify these models so that there is some extra enhancement in the interference part.

3.4 Review of 2HDM

In SM there is only one fundamental scalar and that is the Higgs particle. But there could be more scalars in nature, fundamental or composite. One of the simplest extensions of the scalar sector of the SM is to add a new scalar doublet under the $SU(2)_L$ gauge group to it. This extension of SM is known as the two-Higgs-doublet model (2HDM). In the case of most general scalar potential, the vacuum structure of the 2HDM is very rich. It has 14 parameters and can have CP conserving, CP violating, and charged-violating minima. But we can impose few reasonable conditions such CP is conserved in the scalar sector (so that no mixing takes place between the scalars and the pseudo-scalars), CP is not spontaneously broken, and due to discrete symmetries no quartic term which are odd in either of the doublets are allowed. Now the most general scalar potential which is consistent with all the fore mentioned conditions as well as the SM gauge groups can be written as [29]

$$V(\phi_1,\phi_2) = m_{11}^2\phi_1^\dagger\phi_1 + m_{22}^2\phi_2^\dagger\phi_2 - m_{12}^2(\phi_1^\dagger\phi_2 + h.c) + \frac{\lambda_1^2}{2}(\phi_1^\dagger\phi_1)^2 + \frac{\lambda_2^2}{2}(\phi_2^\dagger\phi_2)^2 \\ +\lambda_3\phi_1^\dagger\phi_1\phi_2^\dagger\phi_2 + \lambda_4\phi_1^\dagger\phi_2\phi_2^\dagger\phi_1 + \frac{\lambda_5^2}{2}[(\phi_1^\dagger\phi_2)^2 + (\phi_2^\dagger\phi_1)^2], \tag{3.4}$$

where all the parameters are taken to be real.

Yukawa sector of 2HDM

The most general Yukawa couplings for quarks in the 2HDM can be written as [29]

$$-\mathcal{L}_{Yukawa} = \bar{Q}_L\eta_1^U U_R\tilde{\phi_1} + \bar{Q}_L\eta_1^D D_R\phi_1 + \bar{Q}_L\eta_2^U U_R\tilde{\phi_2} + \bar{Q}_L\eta_2^D D_R\phi_2 + h.c.\ , \tag{3.5}$$

where η_i with $i = 1,2$ are 3×3 matrices. Similar Yukawa terms can be written for the leptonic sector. Assuming the VEV are real and positive, the mass matrices can be expressed as [29],

$$M^F = \frac{\mathcal{V}_0}{\sqrt{2}}(\eta_1^F\cos\beta + \eta_2^F\sin\beta)\ , \tag{3.6}$$

where F refers to the quark given as $F = U, D$ and $\tan\beta = \frac{v_2}{v_1}$ with the VEV of ϕ_1 and ϕ_2 given as v_1 and v_2 respectively. Denoting

$$\kappa^F = \eta_1^F\cos\beta + \eta_2^F\sin\beta\ , \tag{3.7}$$

and the orthogonal combination as

$$\rho^F = -\eta_1^F\sin\beta + \eta_2^F\cos\beta\ , \tag{3.8}$$

we have in the so-called Higgs basis,

$$\begin{bmatrix} H_1 \\ H_2 \end{bmatrix} = \begin{bmatrix} c_\beta & s_\beta \\ -s_\beta & c_\beta \end{bmatrix} \begin{bmatrix} \Phi_1 \\ \Phi_2 \end{bmatrix}, \tag{3.9}$$

where only H_1 develops an non-zero VEV thus generating fermion masses while the VEV of H_2 is zero i.e $\langle H_1 \rangle = \frac{\mathcal{V}_0}{\sqrt{2}}$ and $\langle H_2 \rangle = 0$ in this basis. In the Higgs basis the doublets can be parametrized by

$$H_1 = \begin{bmatrix} G^+ \\ \frac{1}{\sqrt{2}}(\mathcal{V}_0 + s_{\beta-\alpha} h - c_{\beta-\alpha} H_0 + i G_0) \end{bmatrix}, H_2 = \begin{bmatrix} H^+ \\ \frac{1}{\sqrt{2}}(c_{\beta-\alpha} h + s_{\beta-\alpha} H_0 + i A) \end{bmatrix}, \quad (3.10)$$

where $G^{\pm}$ and G_0 are the Goldstone bosons and $H^{\pm}$ and A are the charged Higgs and the neutral pseudo-scalar eigenstates. Since the LHC discovery of SM-like Higgs in 2012, no significant deviation has been reported of its properties from that of the SM Higgs, so then it is in the decoupling limit where $s_{\beta-\alpha} \to 1$ and the couplings of the lightest scalar in the 2HDM (in the decoupling limit) are nearly identical to that of the SM Higgs.

CP conserving 2HDM with discrete Z_2 symmetry

In 2HDM there are two Higgs doublets Φ_1 and Φ_2, both transforming under the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ as $(1, 2)_{Y=1}$. After SSB both doublets develop vacuum expectation value (VEV) given as $\Phi_i^0 = \frac{v_i}{\sqrt{2}}$ with $v_1 = \mathcal{V}_0 \cos\beta = \mathcal{V}_0 c_\beta$ and $v_2 = \mathcal{V}_0 \sin\beta = \mathcal{V}_0 s_\beta$ where $\mathcal{V}_0 = 246.2$ GeV. We will assume that the scalar potential is real so as not to introduce any extra CP violating terms and other complications. Now if we impose a discrete Z_2 symmetry under which the doublets transform as $\Phi_1 = +\Phi_1$ and $\Phi_2 = -\Phi_2$ then the Z_2 symmetry will restrict the 2HDM to a tree level FCNC free like the SM Higgs which also does not introduce any FCNC at tree level. The 2HDM with Z_2 symmetry can give four different types of very interesting Yukawa interactions although these four types can be regarded as a particular cases of aligned 2HDM or A2HDM without requiring to impose the Z_2 symmetry [46]. In Table 3.2 we have given the four types of 2HDM with which the quark and lepton doublets interact with which of the two Higgs doublets. The Yukawa couplings of the charged Higgs can then be written as [29]

$$\mathcal{L}_{H^{\pm}} = H^+ \bar{U}(V_{CKM} \bar{\rho}^D P_R - \bar{\rho}^U V_{CKM} P_L) D + h.c. \,, \quad (3.11)$$

Model	Type-I	Type-II GeV	l-specific (Type-X)	Flipped (Type-Y)
u_R^i	Φ_2	Φ_2	Φ_2	Φ_2
d_R^i	Φ_2	Φ_1	Φ_2	Φ_1
e_R^i	Φ_2	Φ_1	Φ_1	Φ_2

Table 3.2: The four types 2HDM with Z_2 symmetry which lead to natural flavor conservation (NFC). The superscript i is a generation index. By convention, the ϕ_2 always couple to u_R^i. Taken from reference [29].

where bar over the $\rho^{U,D}$ is to remind that the $\rho^{U,D}$ matrices have been rotated by the same matrices that have diagonalized the mass matrices $\kappa^{U,D}$. Now even in the SM the charged currents can induce flavor changing interactions at tree level but here in the most general case of Yukawa as given above, there will be FCNC at tree level in the neutral current unless $\bar{\rho}^U$ and $\bar{\rho}^D$ are diagonal as well i.e., there will be tree level FCNC unless the matrices that diagonalize the $\kappa^{U,D}$ also diagonalize the $\rho^{U,D}$ simultaneously. Since FCNC are highly constrained, we need to avoid its presence at tree level in any reasonably simple model building. As mentioned before, one of the simplest way to avoid tree level FCNC is to impose the discrete Z_2 symmetry. Under the Z_2 symmetry we have for instance, for $\eta_1^U = \eta_1^D = 0$ (type I 2HDM) and for $\eta_2^D = \eta_1^U = 0$ (type II 2HDM) etc. The charged Higgs interactions in all the four different types of natural flavor conserving 2HDM can be expressed as

$$\mathcal{L}_{H^\pm} = -\left[\frac{\sqrt{2}V_{ud}}{\mathcal{V}_0}\bar{u}(m_u\xi^u P_L + m_d\xi^d P_R)dH^+ + \frac{\sqrt{2}\xi^l m_\tau}{\mathcal{V}_0}\bar{\nu_L} l_R H^+ + h.c.\right] , \quad (3.12)$$

where the values of the $\xi^{u,d,l}$ coupling for a given type of NFC of Table 3.2 is tabulated in Table 3.3.

Model	Type-I	Type-II	Type-X	Flipped (Type-Y)
ξ^u	$\cot\beta$	$\cot\beta$	$\cot\beta$	$\cot\beta$
ξ^d	-$\cot\beta$	$\tan\beta$	-$\cot\beta$	$\tan\beta$
ξ^l	-$\cot\beta$	$\tan\beta$	$\tan\beta$	-$\cot\beta$

Table 3.3: Yukawa couplings of the charged Higgs in the four types of natural flavor conserving 2HDM [29].

3.5 Theoretical Framework

We assume that all the neutrinos are left handed, then the most general effective Hamiltonian that contains all possible four-fermion operators of dimension four for the decay process $b \to c l \nu_l$, where $l = \tau$, μ or e here, is given as [40]

$$\mathcal{H}_{eff} = \frac{4G_F}{\sqrt{2}} V_{cb}[(\delta_{ll} + C^l_{V_L})\mathcal{O}^l_L + C^l_{V_R}\mathcal{O}^l_R - C^l_{S_L}\mathcal{O}^l_{S_L} - C^l_{S_R}\mathcal{O}^l_{S_R} + C^l_T\mathcal{O}^l_T] \ , \quad (3.13)$$

with the operators defined as

$$\mathcal{O}^l_L = (\bar{c}_L\gamma^\mu b_L)(\bar{l}_L\gamma_\mu\nu_{lL}),\ \mathcal{O}^l_R = (\bar{c}_R\gamma^\mu b_R)(\bar{l}_L\gamma_\mu\nu_{lL}),$$

$$\mathcal{O}^l_{S_L} = (\bar{c}_R b_L)(\bar{l}_R\nu_{lL}),\ \mathcal{O}^l_{S_R} = (\bar{c}_L b_R)(\bar{l}_R\nu_{lL}) \text{ and}$$

$$\mathcal{O}^l_T = (\bar{c}_R\sigma^{\mu\nu} b_R)(\bar{l}_L\gamma_{\mu\nu}\nu_{lL}). \quad (3.14)$$

In Eq. (3.13) we have explicitly shown the relative negative sign between effective four fermion operators due to exchange of heavy scalar particles and heavy vector particles. This is due to sign difference between a scalar propagator and a vector propagator[1]. In many analysis the relative sign is implicitly absorbed into the effective coefficients, but

[1]This is why in forces mediated by exchange of scalars, particles carrying same charges attract each other while in forces mediated by exchange of vector particles, particles with same charges repel each other.

if the relative signs between the vector four current operators and the scalar four current operators are explicitly shown, then it will be able to help us rule out few models, where NP is scalar and the real parts of $C^l_{S_L}$ and $C^l_{S_R}$ are dominant (given that we expect NP contribution to be less than the SM contribution). For instance, in 2HDM of type-I and type-II, the effective couplings are real and positive and so in these type of models, the couplings will interfere destructively with that of the SM due to the relative negative sign. Hence, 2HDM of type-I and type-II can only reduce the values. So it is clear from this that the relative sign can actually help us in ruling out all the models of new scalar particles whose effective couplings are non-negative for most parts of the parameter space where NP part is less than the SM part. In this work we will not deal with new vector and tensor terms.

3.6 Modifying Yukawa sector of Flipped 2HDM

In general 2HDM the Yukawa Lagrangian can be written as [29]:

$$\begin{aligned} \mathcal{L}^{2HDM}_{Yukawa} = &- \sum_{f=u,d,l} \frac{m_f}{\mathcal{V}_0} \left(\xi^f_h \bar{f} f h + \xi^f_H \bar{f} f H - i \xi^f_A \bar{f} \gamma_5 f A \right) \\ &- \left[\frac{\sqrt{2} V_{ud}}{\mathcal{V}_0} \bar{u} (m_u \xi^u_A P_L + m_d \xi^d_A P_R) d H^+ + \frac{\sqrt{2} \xi^l_A m_\tau}{\mathcal{V}_0} \bar{\nu_L} l_R H^+ + h.c. \right] , \end{aligned} \tag{3.15}$$

where ξ^f_i with $i = h, H, A$ and $f = u, d, l$ depending on the type of 2HDM being used; see Table 3.3 for the Yukawa coupling in the charged Higgs case, and $\mathcal{V}_0 = 246$ GeV is the vacuum expectation value of the Higgs. But since we require a constructive interference of SM and Charged Higgs contribution to fit all three of $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau \nu_\tau)$, only Lepton specific and Flipped 2HDM can achieve it. But contributions to the interference from just Lepton specific or Flipped 2HDM turn out to be too small to fit the three data simultaneously since the $\tan\beta$ dependence between m_b and m_τ cancels out in these models although they give constructive interference with SM unlike Type-I and Type-II 2HDM. Some additional factor has to be introduced into these models so that it can fit the three data. One simple and straight forward way to achieve this enhancement is if we require that τ lepton is screened from interacting with the full strength to the Higgs VEV v_2 and to all the neutral excitations from Higgs vacuum like the scalars h, H^0 and the pseudo-scalar A^0 in Flipped 2HDM. The details of how such an anomalous interaction can be embedded inside a larger model is outside the scope of the present work and will be assigned to a future work. Here we focus on the phenomenological consequences of such a screening on observed B decay anomalies [27]. Now since τ lepton is screened from seeing the full depth of v_2, its Yukawa coupling must increase so that its mass which is proportional to $Y_{Yukawa} \times \frac{v_2}{\eta}$ now is the same as the observed mass m_τ, which will effectively enhance the Yukawa coupling of τ lepton to the charged Higgs by a factor of η while Yukawa coupling of τ to h, H^0 and A^0 remains same as in usual Flipped 2HDM as the η factors in the neutral scalar Yukawa interaction cancels out with that of the η factor coming from reduced Higgs vacuum $\mathcal{V}_0$. For same analysis carried out in the Lepton Specific 2HDM, see the comments at the end of the chapter. Then the ξ_A^f factors in the charged Higgs interactions in this anomalous Flipped 2HDM with the screening factor η is given by [27] :

$$\xi_A^u = \cot\beta,\ \xi_A^d = \tan\beta \text{ and } \xi_A^{e,\mu} = -\cot\beta \text{ and } \xi_A^\tau = -\eta\cot\beta. \tag{3.16}$$

Finally the charged Higgs Lagrangian contributing to $b \to cl\nu$ in this model can be written in the most general form using Yukawa couplings derived from ξ_A^fs as:

$$\mathcal{L}_{H^\pm}^{Yukawa} = -[V_{cb}\bar{c}(g_s + g_p\gamma_5)bH^+ - \bar{\nu}_l(f_s^l + f_p^l\gamma_5)lH^+ + h.c.]\ , \tag{3.17}$$

where

$$\begin{aligned} g_s &= \frac{(m_b\tan\beta + m_c\cot\beta)}{\sqrt{2}\mathcal{V}_0}, g_p = \frac{(m_b\tan\beta - m_c\cot\beta)}{\sqrt{2}\mathcal{V}_0}\ , \\ f_s^{e,\mu} &= f_p^{e,\mu} = \frac{m_{e,\mu}\cot\beta}{\sqrt{2}\mathcal{V}_0} \text{ and } f_s^\tau = f_p^\tau = \frac{m_\tau\cot\beta}{\sqrt{2}\mathcal{V}_0}\eta\ . \end{aligned} \tag{3.18}$$

Note the relative negative sign between the hadronic current and leptonic current which will ensure constructive interference between SM and charged Higgs contributions. In this model, SM and charged Higgs interfere constructively unlike Type-I and Type-II 2HDM models where they interfere destructively. In this modified Flipped 2HDM, the charged Higgs coupling to τ is enhanced by a factor of η compared to that of a simple Flipped 2HDM.

Then the effective Lagrangian given by charged Higgs exchange between the hadronic and leptonic currents for the $b \to c l \nu_l$ is given as :

$$T_H = +\frac{V_{cb}}{M_H^2}\left[\bar{c}(g_s + g_p\gamma_5)b\bar{l}(f_s^l - f_p^l\gamma_5)\nu_{\bar{l}}\right] . \tag{3.19}$$

And the total effective Lagrangian of SM + charged Higgs for the $b \to c l \nu_{\bar{l}}$ is given as:

$$T_T = \frac{G_F^l V_{cb}}{\sqrt{2}}[\bar{c}\gamma^\mu(1-\gamma_5)b\bar{l}\gamma_\mu(1-\gamma_5)\nu_{\bar{l}}] + \frac{V_{cb}}{M_{H^\pm}^2}[\bar{c}(g_s + g_p\gamma_5)b\bar{l}(f_s^l - f_p^l\gamma_5)\nu_{\bar{l}}] \tag{3.20}$$

where G_F is the Fermi coupling constant and $M_{H^\pm}$ is the mass of the charged Higgs. In what follows we define $s = q^2 = (P_B - P_{D^{(*)}})^2$.

3.7 $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ in the model

In the SM the $B \to D^{(*)}\tau\nu_\tau$ decay occurs via W boson exchange as shown in Fig. 3.2. In SM within the negligible effect due to mass difference between τ, μ and e, the branching fractions of different leptonic final states are expected to be comparable. However, if there exist NP that preferentially impacts heavy fermions, like in Multiple Higgs Doublet Models (MHDM), the effect of these kind of NP will definitely impact $B \to D^{(*)}\tau\nu_l$ differently than $B \to D^{(*)}l\nu_l$, where $l = \mu$ or e, and BABAR, Belle and LHCb report of about 4σ deviation in $R(D^{(*)})$ that may be a hint of a NP that preferentially impacts fermions according to their masses, like 2HDM. In this analysis we examine the implications of the charged Higgs mediated flavor universality violation in a Flipped 2HDM with anomalous charged Higgs coupling to τ lepton as proposed in Section 3.6.

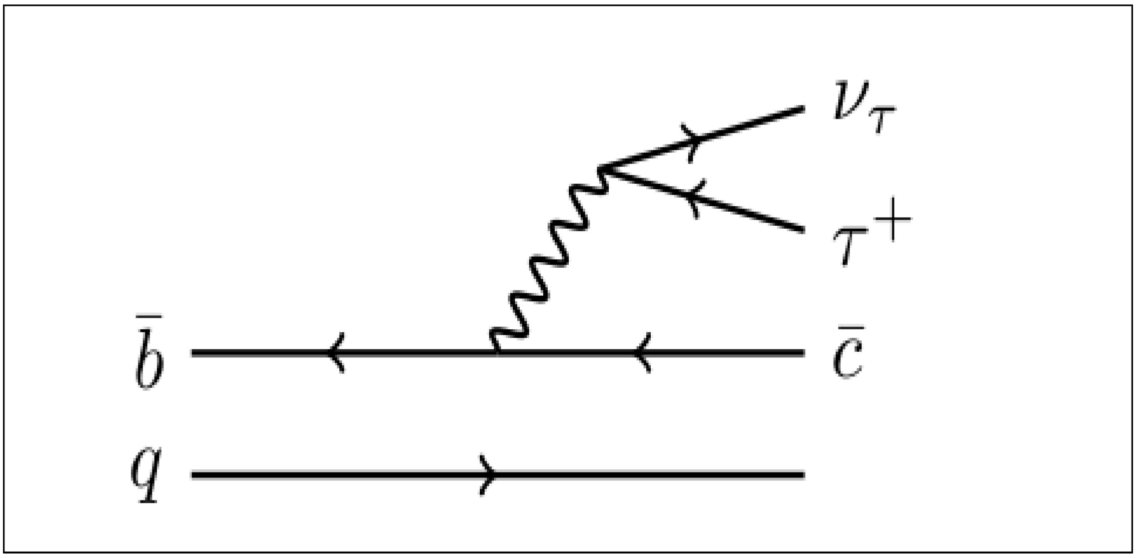

Figure 3.2: Standard-model Feynman diagrams for $B \to D^{(*)}\tau\nu_l$. Similar diagrams can be drawn for $B \to D^{(*)}l\nu_l$, where $l = \mu$ or e but decay rate involving $l = \mu$ or e are not that interesting from the perspective of NP related to Higgs mechanism such as 2HDM etc [48].

$\mathbf{B} \to \mathbf{D}\ l\ \nu_l$

The standard parametrization of the hadronic matrix elements for the vector current is given as

$$\langle D(P_2)|\bar{c}\gamma^\mu(1-\gamma_5)b|B(P_1)\rangle = F_1(s)(P_1+P_2)^\mu + \frac{m_B^2 - m_D^2}{s}[F_0(s) - F_1(s)]q^\mu \ , \quad (3.21)$$

and the standard parametrization of the hadronic matrix elements for the scalar current is given as

$$\langle D(P_2)|\bar{c}b|B(P_1)\rangle = \frac{m_B^2 - m_D^2}{m_b - m_c}F_0 \ , \quad (3.22)$$

where F_0 and F_1 are form factors.

Then we have for the B $\to$ D l ν_l:

$$\frac{d\Gamma_{SM}}{ds} = \frac{G_F^2|V_{cb}|^2}{96\pi^3 m_B^2}\left\{4m_B^2|\vec{P}_D|^2(1+\frac{m_l^2}{2s})|F_1|^2+[m_B^4(1-\frac{m_D^2}{m_B^2})^2\frac{3}{2}\frac{m_l^2}{s}]|F_0|^2\right\}\left(1-\frac{m_l^2}{s}\right)^2|\vec{P}_D|$$

$$\frac{d\Gamma_{MIX}}{ds} = +\frac{G_F}{\sqrt{2}}\frac{1}{M_H^2}\frac{m_l g_s|V_{cb}|^2}{32\pi^3}\left(f_s^l+f_p^l\right)\left[(1-\frac{m_D^2}{m_B^2})\right]\frac{m_B^2-m_D^2}{m_b-m_c}|F_0|^2\left(1-\frac{m_l^2}{s}\right)^2|\vec{P}_D| \quad (3.23)$$

$$\frac{d\Gamma_H}{ds} = \frac{1}{M_H^4}\frac{g_s^2|V_{cb}|^2}{64\pi^3 m_B^2}\left(\frac{m_B^2-m_D^2}{m_b-m_c}\right)^2\left((f_s^l)^2+(f_p^l)^2\right)|F_0|^2 s\left(1-\frac{m_l^2}{s}\right)^2|\vec{P}_D| \quad (3.24)$$

where $|\vec{P}_D| = \frac{\sqrt{s^2+m_B^4+m_D^4-2(sm_B^2+sm_D^2+m_D^2m_B^2)}}{2m_B}$ is the momentum of the D in the B's rest frame and g_s, g_s, f_s^l and f_p^l are taken from Eq. (3.18). In terms of the Babar's parametrization [20] we have

$$F_1 = \frac{\sqrt{m_B m_D}(m_B+m_D)\sqrt{w^2-1}}{2m_B|\vec{P}_D|}V_1 \quad (3.25)$$

$$F_0 = \frac{\sqrt{m_B m_D}(w+1)}{m_B+m_D}S_1 \quad (3.26)$$

where

$$V_1(w) = V(1)[1-8\rho_D^2 z(w)+(51\rho_D^2-10)z(w)^2-(252\rho_D^2-84)z(w)^3] \quad (3.27)$$

$$S_1(w) = V_1(w)\{1+\Delta[-0.019+0.04(w-1)-0.015(w-1)^2]\} \quad (3.28)$$

with $\Delta = 1\pm 1$ and

$$w = \frac{m_B^2+m_D^2-s}{2m_B m_D} \quad (3.29)$$

$$z(w) = \frac{(\sqrt{w+1}-\sqrt{2})}{(\sqrt{w+1}+\sqrt{2})} \quad (3.30)$$

$$\rho_D^2 = 1.186\pm 0.055 \quad (3.31)$$

and common normalization factor $V_1(1)$ cancels in the ratios.
We take

$$R(D) = \frac{\Gamma_T(B\to D\tau\nu_\tau)}{\Gamma_T(B\to Dl\nu_l)} \text{ with } \Gamma_T = \Gamma_{SM}+\Gamma_{MIX}+\Gamma_{H^\pm} \quad (3.32)$$

where l here refers to μ or e.

$\mathbf{B} \to D^* \; l \; \nu_l$

The non-perturbative QCD interactions in the $B \to D^* l\nu_l$ decays can be expressed following the notations of [20], in terms of four independent form factors denoted by V, A_0, A_1, and A_2 as:

$$\langle D^*(p_{D^*}, \varepsilon^*)|\bar{c}\gamma^\mu b|\bar{B}_p\rangle = \frac{iV}{m_B + m_{D^*}} \varepsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu} p_B^\alpha p_{D^*}^\beta \; , \tag{3.33}$$

$$\begin{aligned} \langle D^*(p_{D^*}, \varepsilon^*)|\bar{c}\gamma^\mu\gamma^5 b|\bar{B}_p\rangle = 2m_{D^*} A_0 \frac{\varepsilon^* \cdot q}{s} q^\mu + (m_B + m_{D^*}) A_1 (\varepsilon^* - \frac{\varepsilon^* \cdot q}{s} q^\mu) \\ -A_2 \frac{\varepsilon^* \cdot q}{m_B + m_{D^*}} [(p_B - p_{D^*})^\mu - \frac{m_B^2 - m_{D^*}^2}{s} q^\mu] \; , \end{aligned} \tag{3.34}$$

and we have that the scalar current vanishes due to parity conservation of strong interaction. While using equations of motion, we have the pseudo-scalar current given as: ,

$$\langle D^*(p_{D^*}, \varepsilon^*)|\bar{c}\gamma^5 b|\bar{B}_p\rangle = -\frac{2m_{D^*}}{\bar{m}_b + \bar{m}_c} A_0 \varepsilon^* \cdot q \; . \tag{3.35}$$

Then we have for the B $\to D^* \; l \; \nu_l$:

$$\frac{d\Gamma_{SM}}{ds} = \frac{G_F^2 |V_{cb}|^2 |\vec{P^*}| s}{96\pi^3 m_B^2} \left(1 - \frac{m_\tau^2}{s}\right)^2 \left[(|H_+|^2 + |H_-|^2 + |H_0|^2)(1 + \frac{m_\tau^2}{2s}) + \frac{3m_\tau^2}{2s}|H_s|^2\right] \; , \tag{3.36}$$

$$\frac{d\Gamma_{MIX}}{ds} = +\frac{G_F}{\sqrt{2}} \frac{1}{M_H^2} \frac{m_l g_p |V_{cb}|^2}{8\pi^3} \left(f_s^l + f_p^l\right) \frac{1}{m_b + m_c} |A_0|^2 \left(1 - \frac{m_l^2}{s}\right)^2 |\vec{P}_{D^*}|^3 \; , \tag{3.37}$$

$$\frac{d\Gamma_H}{ds} = \frac{1}{M_H^4} \frac{g_p^2 |V_{cb}|^2}{16\pi^3} \frac{1}{(m_b + m_c)^2} \left((f_s^l)^2 + (f_p^l)^2\right) |A_0|^2 \left(1 - \frac{m_l^2}{s}\right)^2 s |\vec{P}_{D^*}|^3 \; , \tag{3.38}$$

where $|\vec{P}_{D^*}| = \frac{\sqrt{s^2 + m_B^4 + m_{D^*}^4 - 2(s m_B^2 + s m_{D^*}^2 + m_{D^*}^2 m_B^2)}}{2m_B}$ is the momentum of D^* in B's rest frame and g_s, g_s, f_s^l and f_p^l are taken from Eq. (3.18); more details on the helicity amplitudes H_+, H_-, H_0 and H_s and also coordinate frame, kinematics and polarization vectors used in defining the leptonic and hadronic currents are given in the Appendix C. In Babar's parametrization [20] we have

$$H_\pm(s) = (m_B + m_{D^*}) A_1(s) \mp \frac{2m_B}{m_B + m_{D^*}} |\vec{P}_{D^*}| V(s) \; , \tag{3.39}$$

$$H_0(s) = \frac{-1}{2m_{D^*}\sqrt{s}} \left[\frac{4m_B^2|\vec{P}_{D^*}|^2}{m_B + m_{D^*}} A_2(s) - (m_B^2 - m_{D^*}^2 - s)(m_B + m_{D^*}) A_1(s) \right] \ , \quad (3.40)$$

$$H_s(s) = \frac{2m_B|\vec{P}_{D^*}|}{\sqrt{s}} A_0(s) \ , \quad (3.41)$$

where

$$A_1(w) = \frac{w+1}{2} r_{D^*} h_{A_1}(w) \ , \quad (3.42)$$

$$A_0(w) = \frac{R_0(w)}{r_{D^*}} h_{A_1}(w) \ , \quad (3.43)$$

$$A_2(w) = \frac{R_2(w)}{r_{D^*}} h_{A_1}(w) \ , \quad (3.44)$$

$$V(w) = \frac{R_1(w)}{r_{D^*}} h_{A_1}(w) \ , \quad (3.45)$$

with $w = \frac{m_B^2 + m_{D^*}^2 - s}{2m_B m_{D^*}}$ and $r_{D^*} = \frac{2\sqrt{m_B m_{D^*}}}{(m_B + m_{D^*})}$.

The FF are given as:

$$h_{A_1}(w) = h_{A_1}(1)[1 - 8\rho_{D^*}^2 z(w) + (53\rho_{D^*}^2 - 15)z(w)^2 - (231\rho_{D^*}^2 - 91)z(w)^3] \quad (3.46)$$

$$R_0(w) = R_0(1) - 0.11(w-1) + 0.01(w-1)^2 \quad (3.47)$$

$$R_1(w) = R_1(1) - 0.12(w-1) + 0.05(w-1)^2 \quad (3.48)$$

$$R_2(w) = R_2(1) + 0.11(w-1) - 0.06(w-1)^2 \quad (3.49)$$

where $z(w) = \frac{(\sqrt{w+1} - \sqrt{2})}{(\sqrt{w+1} + \sqrt{2})}$ and $\rho_{D^*}^2 = 1.207 \pm 0.028$, $R_0(1) = 1.14 \pm 0.07$, $R_1(1) = 1.401 \pm 0.033$ and $R_2(1) = 0.854 \pm 0.020$ and common normalization factor $h_{A_1}(1)$ cancels in the ratios and we take

$$R(D^*) = \frac{\Gamma_T(B \to D^* \tau \nu_\tau)}{\Gamma_T(B \to D^* l \nu_l)} \text{ with } \Gamma_T = \Gamma_{SM} + \Gamma_{MIX} + \Gamma_{H^\pm} \quad (3.50)$$

where l here refers to μ or e.

$\mathcal{B}\mathbf{r}(\ B^- \to \tau^-\ \nu_\tau)$

In the SM the $B \to \tau\nu_\tau$ decay occurs via W boson exchange as shown in Fig. 3.3 . In presence of NP of charged Higgs type, the total effective Lagrangian of SM + charged Higgs for the $B^- \to \tau^-\nu_\tau$ is given as:

$$T_T = \frac{G_F V_{ub}}{\sqrt{2}}\left[\bar{u}\gamma^\mu(1-\gamma_5)b\bar{\tau}\gamma_\mu(1-\gamma_5)\nu\right] + \frac{V_{ub}}{M_H^2}\left[\bar{u}(g_s' + g_p'\gamma_5)b\bar{\tau}(f_s^\tau - f_p^\tau\gamma_5)\nu\right] \ , \quad (3.51)$$

where $g_s' = \frac{(m_b \tan\beta + m_u \cot\beta)}{\sqrt{3}\mathcal{V}_0}$, $g_p' = \frac{(m_b \tan\beta - m_u \cot\beta)}{\sqrt{2}\mathcal{V}_0}$ and f_s^τ, f_p^τ are taken from Eq. (3.18).

The hadronic matrix elements of pseudo-vector current are given by

$$\langle 0|\bar{u}\gamma^\mu\gamma_5 b|B\rangle = iP_B^\mu f_B \ , \quad (3.52)$$

and hadronic matrix elements of the pseudo-scalar current is given as

$$\langle 0|\bar{u}\gamma_5 b|B\rangle = if_{B0} \ , \quad (3.53)$$

where $M_{H^\pm}$ is the mass of the charged Higgs and $f_{B0} = -\frac{m_B^2}{m_b+m_u} f_B$.

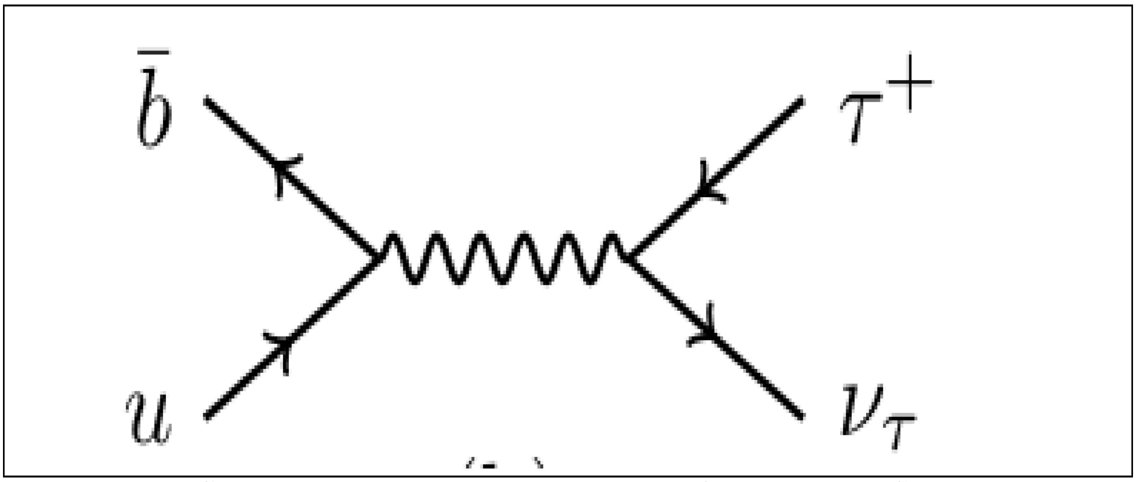

Figure 3.3: Standard-model Feynman diagrams for $B \to \tau\nu_l$. Similar diagrams can be drawn for $B \to l\nu_l$, where $l = \mu$ or e but decay rate involving $l = \mu$ or e are not that interesting from the perspective of NP related to Higgs mechanism such as 2HDM etc [48].

Then we have for the B $\rightarrow \tau\ \nu_\tau$:

$$\mathcal{B}r_{SM}(B \rightarrow \tau\nu_\tau) = \frac{G_F^2 |V_{ub}|^2 m_\tau^2 m_B}{8\pi} f_B^2 (1 - \frac{m_\tau^2}{m_B^2})^2 \tau_B \tag{3.54}$$

$$\mathcal{B}r_{MIX}(B \rightarrow \tau\nu_\tau) = +\frac{G_F}{\sqrt{2}} \frac{1}{M_H^2} \frac{|V_{ub}|^2}{4\pi} g'_p f_B^2 m_\tau \frac{m_B^3}{m_b + m_u} (f_s^\tau + f_p^\tau)(1 - \frac{m_\tau^2}{m_B^2})^2 \tau_B \tag{3.55}$$

$$\mathcal{B}r_{H^\pm}(B \rightarrow \tau\nu_\tau) = \frac{1}{M_H^4} \frac{|V_{ub}|^2}{8\pi} f_B^2 g'^2_p \frac{m_B^5}{(m_b + m_u)^2} ((f_s^\tau)^2 + (f_p^\tau)^2)(1 - \frac{m_\tau^2}{m_B^2})^2 \tau_B \tag{3.56}$$

with

$$\mathcal{B}r_T = \mathcal{B}r_{SM} + \mathcal{B}r_{MIX} + \mathcal{B}r_{H^\pm} \tag{3.57}$$

and τ_B is life time of B meson.

3.8 Results

Although η is an independent parameter but if we require that $\eta = \tan^2\beta$, then it leads to very simple interpretation of the results [27]. If we require $\eta = \tan^2\beta$, then the Yukawa interactions in the hadronic sector of our model has the same form as in the Type-II 2HDM but interaction in the leptonic sector of our model has the same form as in Flipped 2HDM except the τ lepton coupling to the charged Higgs, which has same form as in Type-II 2HDM but with opposite sign. Now there are only two parameters to fit i.e., $\tan\beta$ and $M_\pm$, and a χ^2 analysis with the $R(D^{(*)})$ and $\mathcal{B}$r($B^- \rightarrow \tau^-\ \nu_\tau$) as data points to find the best fits for the parameters $\tan\beta$ and $M_\pm$ gives $\chi^2_{min} = 10.95$ for $M_\pm > 600$ GeV. This $M_\pm > 600$ GeV constraint is from $B \rightarrow X_s\gamma$ and latest estimate on the lower bound of $M_\pm$ from $B \rightarrow X_s\gamma$ is given in [30]. We have tabulated for two different values of the parameters that fits at same accuracy in the table below [27] :

S.no	$\tan\beta$	$M_{\pm}$ GeV	$R(D)_{Th}$	$R(D^*)_{Th}$	$Br_{Th}(B \to \tau\nu)$
1	69.97	700	0.348	0.255	1.29×10^{-4}
2	99.95	1000	0.348	0.255	1.29×10^{-4}

Table 3.4: $\chi^2_{min} = 10.95$ and we have restricted the $\tan\beta$ in the range of $100 > \tan\beta > 1$ [27].

As the data in the table above and combined errors from experiments given in Eqs(3.1) and Eqs(3.3) shows, in the range 1 TeV $\geq M_{\pm} \geq$ 600 GeV and $100 > \tan\beta > 1$, we have [27] :

$$R(D)_{Th} = 0.348 \pm 0.16 \ , \tag{3.58}$$

$$R(D^*)_{Th} = 0.255 \pm 0.07 \ , \tag{3.59}$$

and

$$Br_{Th}(B \to \tau\nu) = (1.29 \pm 0.89) \times 10^{-4} \ , \tag{3.60}$$

compared to the combined [Babar,Belle,LHCb] [23] experimental values:

$$R(D)_{EXP} = 0.388 \pm 0.047 \ , \tag{3.61}$$

$$R(D^*)_{EXP} = 0.321 \pm 0.021 \ , \tag{3.62}$$

and

$$Br_{EXP}(B \to \tau\nu) = (1.14 \pm 0.27) \times 10^{-4} \ . \tag{3.63}$$

3.9 Summary

Babar, Belle and recently LHCb has reported an excess in the measurements of $R(D^*)$, $R(D)$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ larger than expected from SM, a possible signature of lepton flavor universality violating NP. In this work we have analyzed the implications for these decay modes in a Flipped 2HDM with enhanced Yukawa coupling of $H^{\pm}$ to τ lepton [27]. By adding theoretical and experimental errors in quadrature from Eqs. (3.58,3.59,3.60) and Eqs. (3.61,3.62,3.63), we conclude that our phenomenological model can give results in agreement within 1σ deviation for the combination of $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_{\tau})$ compared to about 4σ deviation from SM from the latest combined [Babar,Belle,LHCb] experimental data for these observables.

The same results can be achieved if b quark replaces the τ lepton in a Lepton Specific 2HDM. In that case Yukawa coupling of the leptonic sector will be same as Type-II 2HDM and Yukawa coupling of the quark sector will be same as Lepton Specific 2HDM [29] except the b quark which will have the effective Yukawa coupling same as in Type-II with opposite sign. In that case charged Higgs mass may be allowed to be lower than 600 GeV; for a recent work on muon $g-2$ in Lepton Specific 2HDM and related bounds see [30, 31].

From the form of the Yukawa couplings it is expected that if we require $\eta = -1$ for the b quark or τ lepton in the 2HDM-II, then also it will fit the three data at about same accuracy as above model of Flipped/Lepton-Specific 2HDM. This is interesting in a sense that it will be like 2HDM-II but with a wrong sign in either b quark or τ lepton Yukawa coupling, an anomalous SUSY.

Chapter 4

CP Violation in Hadronic τ Decays

4.1 Introduction

The violation of combined symmetries of charge conjugation and parity (CP) has been of interest and searched for a long time as a possible source of the origin of the matter anti-matter asymmetry in the universe. So far most of the efforts to search for CP violation have been in the hadronic sector. In SM, the only source of CP violation is the one phase in the Kobayashi Maskawa (KM) quark mixing matrix. While the Kobayashi Maskawa ansatz for CP violation within the Standard Model [1] in the quark sector has been clearly verified by the plethora of data from the B factories, this is unable to account for the observed baryon asymmetry of the Universe. The observation of neutrino oscillations makes it important to re-examine the question of CP violation in the leptonic sector. If CP violation is observed in the leptonic sector then it is a clear indication of NP beyond SM. However CP violation in the leptonic sector has not been observed yet. CP violation in the neutrino sector can be explored by measuring the asymmetry of the oscillation probabilities in for instance $\nu_\tau \to \nu_\mu$ and $\bar{\nu}_\tau \to \bar{\nu}_\mu$. This kind of CP violation in the neutrino oscillation requires lepton flavor violation as well as non degeneracy of the masses of the charged leptons and neutrinos. Another kind of CP violation in the leptonic sector can be explored using τ leptons produced in the B factories and Super flavor factories. The study of CP violation in τ decays has always been of much interest for beyond the Standard Model studies in the past two decades. Apart from the CP phases that may arise in the neutrino mixing matrix, the decays of the τ lepton may allow us to explore nonstandard CP-violating interactions. Various experimental groups have been involved in exploring CP violation in τ decays in the last decade or more. In 2002, the CLEO collaboration [2], and more recently the Belle Collaboration [3], studied the model dependent angular distribution of the decay products in $\tau \to K_s \pi \nu_\tau$ in search of CP violation coming from charged scalar type of NP; however neither study revealed any CP asymmetry. The BABAR collaboration [20] for the first time reported a sign anomaly in the integrated decay rate CP asymmetry $A_{cp}(\tau \to K_s \pi \nu_\tau)$ at about 2.8σ deviation from the SM prediction originating in the $K^0 - \bar{K}^0$ oscillation [12][37]. More details about the observed sign anomaly in the integrated decay rate CP asymmetry $A_{cp}(\tau \to K_s \pi \nu_\tau)$ in tensor current NP will be presented in Chapter 5.

4.2 CP Violation in The Hadronic τ Decays

In this section we will review the CP violation formulation given in Reference [7] for the scalar and vector type of NP assuming neutrinos are purely left handed and enumerate the complete parameter set for CP violating NP in hadronic τ decays. For the CP violating NP coming from tensor type of NP see Chapter 5. Assuming neutrinos are purely left handed fermions, we can write down the effective Hamiltonian contributing to the τ decays as [7]:

$$\begin{aligned}\mathcal{H}_{eff} = V_{ud}\frac{G_F}{\sqrt{2}}\left[\bar{\nu}\gamma_\alpha(1-\gamma_5)\tau\right]&\Bigg\{[(1+\chi_V^d)g^{\alpha\beta} + \frac{Q^\alpha Q^\beta}{m_\tau(m_u - m_d)}\eta_S^d]\bar{d}\gamma_\beta u \\ &- \left[(1+\chi_A^d)g^{\alpha\beta} + \frac{Q^\alpha Q^\beta}{m_\tau(m_u + m_d)}\eta_P^d\right]\bar{d}\gamma_\beta\gamma_5 u\Bigg\} ,\end{aligned} \tag{4.1}$$

where Q^μ is the total hadronic momentum. Similar terms can be written with the d quark replaced by the s quark. As we can deduce from the above formula and a similar one for the d quark replaced by s quark, CP violation coming from NP of Scalar and/or Vector type of new particle exchange is completely described by eight parameters of $\chi_V^{d,s}$, $\chi_A^{d,s}$, $\eta_S^{d,s}$ and $\eta_P^{d,s}$, assuming all neutrinos are left handed. Also if CP violation is to come from exchange of new scalar and/or vector particles then it necessarily has to arise from the complex phase contained in at least one of the eight parameters above. With addition of two more parameters coming from NP which can generate effective tensor type of coupling, we have total of 10 complex parameters to pin down CP violating NP in hadronic τ decays.

Two-meson Final States

In general, CP violation from vector type NP will be observable only if both vector current and axial vector currents contribute to the same final states [6, 7]. Since in two pseudo scalar meson final states only vector current can contribute due to parity conservation of strong interaction, and in one vector/axial-vector meson with one pseudo scalar meson final states like $(\tau \to a_1 \pi \nu_\tau)/(\tau \to \rho \pi \nu_\tau)/(\tau \to \omega \pi \nu_\tau)$, which are eigenstates of G parity, only either vector current or the axial-vector current will contribute due to G parity conservation [15] of strong interaction, so vector type of NP can contribute in general to CP violation in three or more pseudo scalar meson final states but not in two pseudo scalar meson final states and one vector/axial-vector meson with one pseudo scalar meson final states which are eigenstates of G parity like $(\tau \to a_1 \pi \nu_\tau)/(\tau \to \rho \pi \nu_\tau)/(\tau \to \omega \pi \nu_\tau)$. So in this section for the two meson final states we will consider CP violation coming from scalar type of NP only and details on CP violation coming from tensor type of NP will be given in Chapter 5. For CP violation coming from scalar type of NP in two meson final states, we are basically probing the parameters $\eta_S^{d,s}$ and $\eta_P^{d,s}$ where in the final states of two pseudo scalar mesons and in the final states of one axial-vector and one pseudo scalar meson, we can only probe $\eta_S^{d,s}$ as only vector and scalar currents will contribute here. Now to probe η_S^d we can either use the two pion final state as was done by CLEO [11] but as argued in Reference [15], CP violation coming from $\pi\pi$ final states from a spin zero charged scalar type of NP is suppressed at least to the level of Iso-spin breaking. So the next best option is either ($\tau \to a_1\pi\nu_\tau$) or ($\tau \to 4\pi\nu_\tau$), but form factors of the four pion final state are not calculated accurately and theoretical efforts on that front are on going so the best option as of now is ($\tau \to a_1\pi\nu_\tau$); here also the form factors are not fitted yet for the ($\tau \to a_1\pi\nu_\tau$) as far as the author knows. The two pseudo scalar meson final state with one kaon $\tau \to K\pi\nu_\tau$ will be able to probe η_S^s parameter, and the final states of $\tau \to 3\pi\nu_\tau$ and $\tau \to K\pi\pi\nu_\tau$ will be able to probe η_P^d and η_P^s respectively as well as CP violating parameters coming from vector type NP i.e., $\chi_V^{d,s}$ and $\chi_A^{d,s}$. In what follows we will use s and Q^2 synonymously unless explicitly stated otherwise.

Two Pseudoscalar Meson Final States: $\tau \to P_1P_2\nu_\tau$

After integration over the unobservable neutrino direction, the differential decay rate in the hadronic rest frame is given as [7]

$$d\Gamma(\tau \to P_1P_2\nu_\tau) = \left\{\bar{L}_BW_B + \bar{L}_{SA}W_{SA} + \bar{L}_{SF}W_{SF} + \bar{L}_{SG}W_{SG}\right\} \times \frac{G_F^2V_{u(d/s)}^2}{2m_\tau}\frac{1}{(4\pi)^3}\frac{(m_\tau^2-s)^2}{m_\tau^2}|\vec{q}_1|\frac{ds}{\sqrt{s}}\frac{d\cos\theta}{2}\frac{d\alpha}{2\pi}\frac{d\cos\beta}{2}\ , \tag{4.2}$$

where

$$W_B = 4(\vec{q}_1)^2|F_V|^2\ , \tag{4.3}$$

$$W_{SA} = s|\bar{F}_S|^2\ , \tag{4.4}$$

$$W_{SF} = 4|\vec{q}_1|\sqrt{s}Re[F_V\bar{F}_S^*]\ , \tag{4.5}$$

$$W_{SG} = -4|\vec{q}_1|\sqrt{s}Im[F_V\bar{F}_S^*] \tag{4.6}$$

with F_S and F_V being the scalar and vector form factors respectively and $|\vec{q}_1| = q_1^z$ is the momentum of the P_1 in the hadronic rest frame and is given as

$$|\vec{q}_1| = \frac{1}{2\sqrt{s}}((s-m_1^2-m_2^2)^2 - 4m_1^2m_2^2)^{1/2}\ , \tag{4.7}$$

with

$$\bar{F}_S = (1 + \frac{s}{m_\tau(m_u - m_d)}\eta_S^{(d/s)})F_S \ . \tag{4.8}$$

If the τ direction in the hadronic rest frame is known i.e., $\psi \to 0$, and $P \neq 0$, then we have

$$\bar{L}_B = K_1 \sin^2\beta + K_2 - K_4 \sin 2\beta \cos\alpha \ , \tag{4.9}$$

$$\bar{L}_{SA} = K_2 \ , \tag{4.10}$$

$$\bar{L}_{SF} = -K_2 \cos\beta - K_4 \sin\beta \cos\alpha \ , \tag{4.11}$$

$$\bar{L}_{SG} = K_5 \sin\beta \sin\alpha \ . \tag{4.12}$$

with

$$K_1 = 1 - \gamma_{VA} P \cos\theta - (m_\tau^2/s)(1 + \gamma_{VA} P \cos\theta) \ , \tag{4.13}$$

$$K_2 = (m_\tau^2/s)(1 + \gamma_{VA} P \cos\theta) \ , \tag{4.14}$$

$$K_3 = \gamma_{VA} - P \cos\theta \ , \tag{4.15}$$

$$K_4 = \sqrt{(m_\tau^2/s)}\gamma_{VA} P \sin\theta \ , \tag{4.16}$$

$$K_5 = \sqrt{(m_\tau^2/s)} P \sin\theta \ , \tag{4.17}$$

$$\bar{K}_1 = K_1(3\cos^2\psi - 1)/2 - 3/2 K_4 \sin 2\psi \ , \tag{4.18}$$

$$\bar{K}_2 = K_2 \cos\psi + K_4 \sin\psi \ , \tag{4.19}$$

$$\bar{K}_3 = K_3 \cos\psi - K_5 \sin\psi \ , \tag{4.20}$$

substituting $\psi \to 0$ in the definitions of the $\bar{K}$'s above and given that the angle θ is the angle between direction of flight of the τ in the laboratory frame and direction of the hadrons as seen in the τ rest frame and the definitions of angles α, β and ψ are given in Fig. 4.1. If the τ direction in the hadronic rest frame is not known then the azimuthal angle α is not observable and has to be integrated out. When τ pairs are produced in symmetric colliders, we can express the angles θ and ψ in terms of s and x as given in Reference [7]

$$\cos\theta = \frac{2xm_\tau^2 - m_\tau^2 - s}{(m_\tau^2 - s)\sqrt{1 - m_\tau^2/E_{Beam}^2}}\ , \tag{4.21}$$

and

$$\cos\psi = \frac{x(m_\tau^2 + s) - 2s}{(m_\tau^2 - s)\sqrt{x^2 - s/E_{Beam}^2}}\ , \tag{4.22}$$

where $E_{beam}^2 = E_\tau^2$ and $x = 2E_h/\sqrt{s}$ with E_h being the energy of the hadrons in the laboratory frame. From Eqs. (4.5) and (4.6) we can see that the only relevant term for the CP violation coming from scalar type NP in two pseudo scalar meson final states are W_{SF} and W_{SG} and if the direction of the τ is not measured, then $\bar{L}_{SG} = 0$ and so W_{SG} is not measurable and only W_{SF} can be measured in that case.

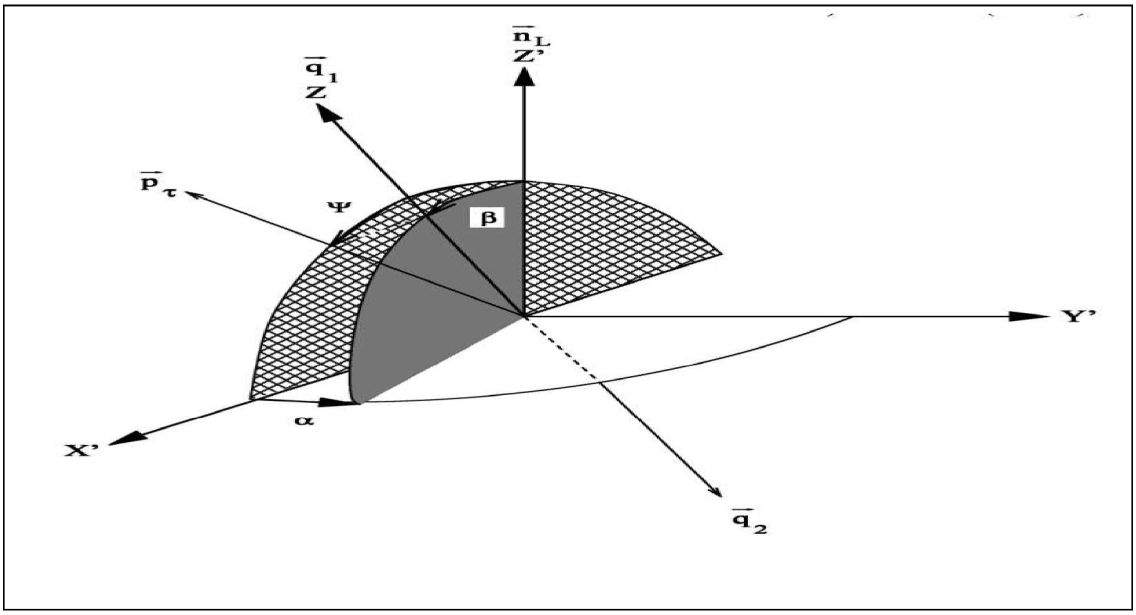

Figure 4.1: Definitions of the angles α, β and ψ [15].

Vector/Axial-Vector and Pseudo-scalar Meson Final States:

$\tau \to V_1 P_2 \nu_\tau$

In general the hadronic current J^μ in the case of the final states of a vector/axial-vector and a pseudo scalar meson like $(\tau \to a_1\pi\nu_\tau)/(\tau \to \rho\pi\nu_\tau)/(\tau \to \omega\pi\nu_\tau)$ can be expressed by four form factors given as [16]:

$$\begin{aligned} J^\mu = F_1(Q^2)\left(Q^2\epsilon_1^\mu - \epsilon_1\cdot q_2 Q^\mu\right) + F_2(Q^2)\epsilon_1\cdot q_2\left(q_1^\mu - q_2^\mu - Q^\mu\frac{Q\cdot(q_1-q_2)}{Q^2}\right) \\ +iF_3(Q^2)\epsilon^{\mu\alpha\beta\gamma}\epsilon_{1\alpha}q_{1\beta}q_{2\gamma} + F_4(Q^2)\epsilon_1\cdot q_2 Q^\mu \ , \end{aligned} \tag{4.23}$$

where ϵ_1 is the polarization vector of the vector/axial-vector meson which satisfies $\epsilon_1\cdot q_1 = 0$ and $\sum_\gamma \epsilon_1^\mu(\gamma)\epsilon_1^{*\nu}(\gamma) = -g^{\mu\nu} + \frac{q_1^\mu q_1^\nu}{m_1^2}$.
After integration over the unobservable neutrino direction and unknown τ direction α, the differential decay rate in the hadronic rest frame is given as [16]

$$\begin{aligned} d\Gamma(\tau \to V_1P_2\nu_\tau) = \{L_AW_A + L_BW_B + L_EW_E + L_{SA}W_{SA} + L_{SF}W_{SF} + L_{SG}W_{SG}\} \times \\ \frac{G_F^2 V_{u(d/s)}^2}{4m_\tau}\frac{1}{(4\pi)^3}\frac{(m_\tau^2 - s)}{m_\tau^2}|\vec{q}_1|\frac{ds}{\sqrt{s}}\frac{d\cos\theta}{2}\frac{d\cos\beta}{2} \end{aligned} \tag{4.24}$$

where

$$L_A = (g_A^2 + g_V^2)(m_\tau^2 - s)[2/3K_1 + K_2 + 1/3\bar{K}_1(3\cos^2\beta - 1)/2] \ , \tag{4.25}$$

$$L_B = (g_A^2 + g_V^2)(m_\tau^2 - s)[2/3K_1 + K_2 - 2/3\bar{K}_1(3\cos^2\beta - 1)/2] \ , \tag{4.26}$$

$$L_E = (g_A^2 + g_V^2)(m_\tau^2 - s)\bar{K}_3\cos\beta \ , \tag{4.27}$$

$$L_{SA} = (g_A^2 + g_V^2)(m_\tau^2 - s)K_2 \ , \tag{4.28}$$

$$L_{SF} = -(g_A^2 + g_V^2)(m_\tau^2 - s)\bar{K}_2\cos\beta \ , \tag{4.29}$$

$$L_{SG} = 0 \ , \tag{4.30}$$

where V_1 is a vector or a axial vector meson and $\gamma_{VA} = \frac{2g_Vg_A}{g_V^2+g_A^2} = 1$ in SM, the angle θ is the angle between direction of flight of the τ in the laboratory frame and direction of the hadrons as seen in the τ rest frame, the angle ψ is the angle between direction of flight of the τ and direction of the laboratory as seen in hadronic rest frame; for more details see Fig. 4.1 on page 68. The values of the K_i with i = 1,2,3,4,5 are the same as given in Section 4.2. When τ pairs are produced in colliders we can express the angles θ and ψ in terms of s and x as given in Eqs. (4.21), (4.22). Now for the hadronic matrix element part, we have

$$W_A = 2s(|\vec{q}_1|^2|F_3|^2 + s|F_1|^2) \ , \tag{4.31}$$

$$W_B = s/m_1^2(sE_1^2|F_1|^2 + 4\sqrt{s}E_1|\vec{q}_1|^2 Re(F_1F_2^*) + 4|\vec{q}_1|^4|F_2|^2) \ , \tag{4.32}$$

$$W_E = 4(\sqrt{s})^3|\vec{q}_1|Re(F_1F_3^*) \ , \tag{4.33}$$

$$W_{SA} = \frac{s^2|\vec{q}_1|^2}{m_1^2}|\bar{F}_4|^2 \ , \tag{4.34}$$

$$W_{SF} = \frac{s^{3/2}|\vec{q}_1|}{m_1^2}(2Re(F_1\bar{F}_4^*)(\sqrt{s}E_1) + 4Re(F_2\bar{F}_4^*)|\vec{q}_1|^2) \ , \tag{4.35}$$

$$W_{SG} = \frac{s^{3/2}|\vec{q}_1|}{m_1^2}(2Im(F_1\bar{F}_4^*)(\sqrt{s}E_1) + 4Im(F_2\bar{F}_4^*)|\vec{q}_1|^2) \ . \tag{4.36}$$

Eq. (4.33) has contribution from both vector (F_3) as well as axial vector (F_1) form factors which therefore contain information about CP violation from vector type NP in general but in the case of $(\tau \to a_1\pi\nu_\tau)/(\tau \to \rho\pi\nu_\tau)/(\tau \to \omega\pi\nu_\tau)$ where only vector or axial vector form factors contributes due to G parity conservation, we cannot extract CP violation information from new vector particle exchanges. Now with $\bar{F}_4 = (1+\frac{s}{m_\tau(m_u-m_d)}\eta^d_{S/P})F_4$, we can extract CP violation information from new scalar particle exchanges from W_{SF} and W_{SG} in general but since $L_{SG} = 0$ from Eq. (4.30), CP violating information contained in the W_{SG} is also lost. F_4 is zero in the approximation where the quark masses can be neglected and therefore in the SM, contribution from F_4 is expected to be small. So a significant contribution from F_4 can only come from a new scalar particle in the form of $\frac{s}{m_\tau(m_u-m_d)}\eta^d_{S/P}F_4$.

Optimal CP Observable in $\tau \to a_1 \pi \nu_\tau$

Since W_{SG} is not observable due to the fact that $L_{SG} = 0$, we have only W_{SF} to probe CP violation from scalar type of NP. Now from the functional form of the W_{SF} from Eq. (4.35) we see that only F_1 and F_2 will contribute to CP observables. From G parity only $a_1 \pi \nu_\tau$ final state can contribute substantially to F_1 and F_2, and also as mentioned before $\tau \to a_1 \pi \nu_\tau$ is the best option to pin down the η_S^d parameter. In Reference [15] it has been argued that if the differential decay width is weighted with $\cos\beta$ and then integrated over the $\cos\beta$, we get a CP observable in $\tau \to a_1 \pi \nu_\tau$ where polarization of the τ may enhance the observable value given as

$$\begin{aligned} A_{CP}^{\langle \cos\beta \rangle}(\tau \to a_1 \pi \nu_\tau) &= \frac{1}{\Gamma_{Total}} \int \left[\frac{d\Gamma}{dsd\cos\beta} - \frac{d\bar{\Gamma}}{dsd\cos\beta} \right] \cos\beta \, ds \, d\cos\beta \\ &\approx \frac{V_{ud}^2}{2m_\tau^6} \int (m_\tau^2 - s)^2/(s)^{3/2} |q_1^z| (W_{SF} - \bar{W_{SF}}) \rho(s) ds \ , \end{aligned} \tag{4.37}$$

where

$$\rho(s) = \int \left(\cos\psi + P\cos\theta\cos\psi + P\frac{\sqrt{s}}{m_\tau}\sin\theta\sin\psi \right) \frac{d\cos\theta}{2} \ , \tag{4.38}$$

with P being the polarization of the τ lepton. Here will also show by the optimal observable method that polarized τ are more sensitive to CP violation than non-polarized τ observables with few assumptions made about the form factors since as of now the form factors F_1 and F_2 have not been fitted properly, and also all fits take $F_4 = 0$ which is valid only in the limit of vanishing quark masses.

One of the simplest methods in searching for CP violation, especially in model dependent CP violation searches, is the Optimal observable method. The optimal CP observable is defined as the CP-odd observable that has the smallest associated statistical error. In a given model, the interaction Lagrangian can be separated into a CP-odd P_{odd} and a CP-even P_{even} part; the optimal observable is defined as $\xi = P_{odd}/P_{even}$. As we can see from the definition of ξ, in order to construct it we need to know the explicit functional forms of the CP-odd and CP-even parts of the amplitude in terms of the experimentally measured parameters of the decay. Explicit functional form can be given only in a given model so choice of ξ is model dependent. Now if we weight the differential width with CP-odd observable ξ and then integrate over the $\cos\beta$ and $\cos\theta$, a value different from zero is a clear indication of CP violation. As argued earlier, only from W_{SF} we will be able to extract CP violation information, so the only source of CP-odd observable is W_{SF}. Since F_4 has not been fitted properly till now we will take it to be constant such that $\frac{1}{(m_u-m_d)}F_4 \approx 10^2$ GeV^{-1}. CP-odd observables are the terms which contain the imaginary part of the η_S^d and the CP-even part are the terms which are independent of η_S^d or terms containing the real part of the η_S^d. Since only W_{SF} can contribute to the CP-odd observable, we have from the form of the W_{SF} from Eqs(4.35) and taking the limit $\mathrm{Re}(\eta_S^d) = 0$, $\mathrm{Im}(\eta_S^d) = 1$ and $\frac{1}{(m_u-m_d)}F_4 \approx 10^2$,

$$P_{odd} = \frac{s^{3/2}|\vec{q}_1|}{m_1^2}\left(-2Im(F_1)\frac{s}{m_\tau}(\sqrt{s}E_1) - 4Im(F_2)\frac{s}{m_\tau}|\vec{q}_1|^2\right) \times 10^2 \ . \tag{4.39}$$

P_{even} is the same as $d\Gamma(\tau \to a_1\pi\nu_\tau)$ from Eq. (4.24) except without the $L_{SF}W_{SF}$ term in that equation. Since form factors are not fitted properly for the $\tau \to a_1\pi\nu_\tau$ we will take the form factors F_1 and F_2 having the same form as for the F_3 which has been used in [15] given as

$$F_{1,2} = \frac{g_{\rho\omega\pi}}{\gamma_\rho}\left[\frac{m_\rho^2}{s - m_\rho^2 + I\sqrt{s}\Gamma_\rho} + A_1\frac{m_{\rho'}^2}{s - m_{\rho'}^2 + I\sqrt{s}\Gamma_{\rho'}}\right] \ , \tag{4.40}$$

where m_ρ, Γ_ρ are taken from PDG values and we take $g_{\rho\omega\pi} = (16.1 \pm 0.6) \times 10^3$ GeV^{-1} as given in model 2 of Table I in CLEO (2000) [17] that has been fitted for the $\tau \to \rho\pi\nu_\tau$ and $\gamma_\rho = 4.95$ and $A_1 = -0.24$. These values are actually from the fits of $\tau \to \rho\pi\nu_\tau$, but here we are using these values only to demonstrate the possible enhancement of the CP violation from scalar type of NP if we use polarized τ over unpolarized τ. Also we take $m_{\rho'} = 1.53$ GeV and

$$\begin{aligned}\Gamma_{\rho'} = \Gamma_{\rho'0}\Bigg[&\mathcal{B}r(\rho' \to \pi\pi)\frac{M_{\rho'}}{|Q|}\left(\frac{p_\pi(s)}{p_\pi(M_{\rho'}^2)}\right)^3 + \mathcal{B}r(\rho' \to \omega\pi)\frac{M_{\rho'}}{|Q|}\left(\frac{p_\omega(s)}{p_\omega(M_{\rho'}^2)}\right)^3 \\ &+ \mathcal{B}r(\rho' \to a_1\pi)\frac{M_{\rho'}}{|Q|}\left(\frac{p_{a_1}(s)}{p_{a_1}(M_{\rho'}^2)}\right)\Bigg] \ ,\end{aligned} \tag{4.41}$$

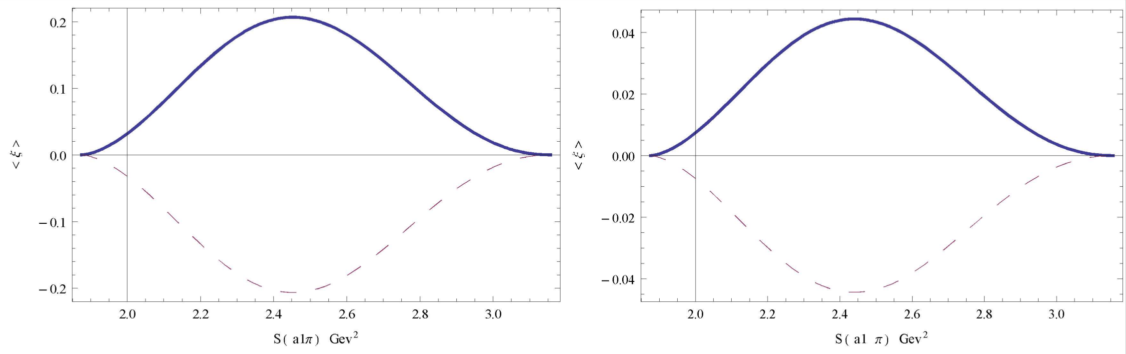

Figure 4.2: Plots of $\langle\xi\rangle(Q^2)$ as a function of $S = Q^2$ for the Polarized τ (left figure) and Unpolarized τ (right figure) where $F_2 = 0$ and F_1 taken from Eq. (4.40).

with $\Gamma_{\rho'0} = 0.43 \pm 0.3$ GeV and all other branching fractions are taken from PDG values except $\mathcal{B}r(\rho' \to a_1\pi)$ which is fixed from $R_{a_1\pi} = \mathcal{B}r(\tau \to a_1\pi\nu_\tau)/\mathcal{B}r(\tau \to 4\pi\nu_\tau)$ with $R_{a_1\pi} = 0.43$. So then we have

$$\langle\xi\rangle(Q^2) = 1/Br(\tau \to 4\pi\nu) \int \frac{P^2_{odd}}{P_{even}} dcos\beta \, d\cos\theta \ , \tag{4.42}$$

where we have normalized the rates by the branching fraction of Br($\tau \to 4\pi\nu$), and assumed only one of the form factors F_1 or F_2 is non zero. Figs. 4.2 and 4.3 show the plots of $\langle\xi\rangle(Q^2)$ as a function of Q^2 for polarized τ and unpolarized τ.

As we can see from Figs. 4.2 and Figure 4.3, the optimal observable $\langle\xi\rangle(Q^2)$ in the case of polarized τ is almost an order of magnitude larger than that of the optimal observable $\langle\xi\rangle(Q^2)$ in the case of unpolarized τ. This makes a good case for using polarized τ in the future whether in the form of correlation measurements or direct polarization.

Three Pseudoscalar Meson Final States

In the semi-leptonic τ decays into three pseudo-scalar meson final states

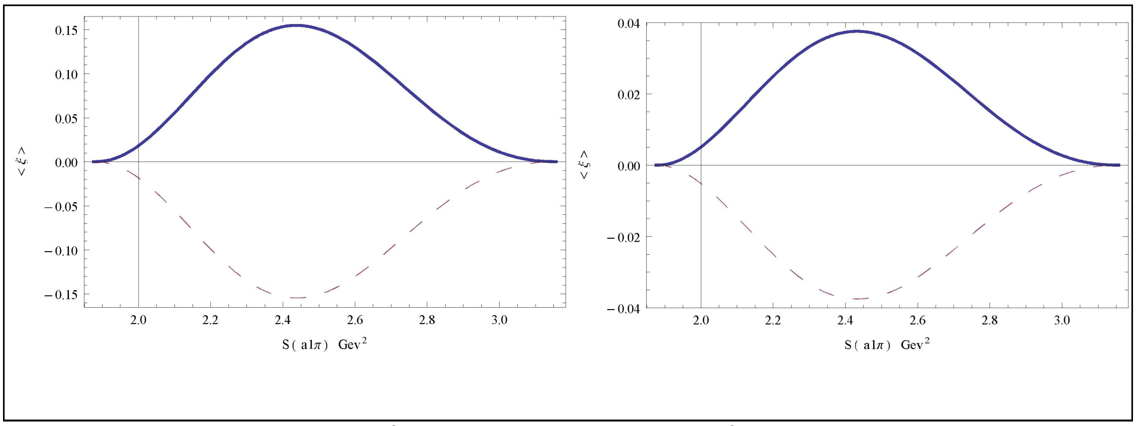

Figure 4.3: Plots of $\langle\xi\rangle\,(Q^2)$ as a function of $S = Q^2$ for the Polarized τ(left one) and Unpolarized τ(right one) where $F_1 = 0$ and F_2 taken from Eqs(4.40).

$$\tau(l,s) \to \nu_\tau(l',s') + h_1(q_1,m_1) + h_2(q_2,m_2) + h_3(q_3,m_3) \tag{4.43}$$

we have [18]

$$d\Gamma = \frac{G_F^2}{4m_\tau} V_{u(d/s)}^2 \{L_{\mu\nu} H^{\mu\nu}\} \times \frac{1}{(2\pi)^5} \frac{1}{64} \frac{(m_\tau^2 - Q^2)}{m_\tau^2} \frac{dQ^2}{Q^2} ds_1 ds_2 \times \frac{d\alpha}{2\pi} \frac{d\gamma}{2\pi} \frac{d\cos\beta}{2} \frac{d\cos\theta}{2}\ , \tag{4.44}$$

where

$$L_{\mu\nu} = \{L^1_{\mu\nu} + L^2_{\mu\nu} - (L^3_{\mu\nu} + L^4_{\mu\nu})\}\ . \tag{4.45}$$

With $\{a,b\}_{\mu\nu} = (a_\mu b_\nu + b_\mu a_\nu - ab g_{\mu\nu})$, we have $L^1_{\mu\nu} = 4\{l,l'\}_{\mu\nu}$, $L^2_{\mu\nu} = -4im_\tau \varepsilon_{\alpha\beta\mu\nu} l'^\alpha s^\beta$, $L^3_{\mu\nu} = -4i\varepsilon_{\alpha\beta\mu\nu} l'^\alpha l^\beta$ and $L^4_{\mu\nu} = 4m_\tau\{S,l'\}_{\mu\nu}$ where S denotes the polarization (P) of the τ satisfying $l_\mu S^\mu = 0$ and $S_\mu S^\mu = -P^2$. The hadronic matrix element part is given as $H^{\mu\nu} = J^\mu J^{\nu\dagger}$ with $J^\mu = \langle h_1(q_1)h_2(q_2)h_3(q_3)|J^\mu_A + J^\mu_V|0\rangle$. The decay rate can be most easily analyzed in the hadronic rest frame $q_1 + q_2 + q_3 = 0$ with introduction of two coordinate frames $S' = Ox'y'z'$ and $S = Oxyz$ shown in Fig. 4.4.

The S' coordinate system allows a simple description of the τ momentum and its spin. The Oz' is pointing in the direction of the laboratory (n_L) viewed from the hadronic rest frame ($n_L = -n_Q$ where n_Q is the direction of the hadrons viewed from the laboratory

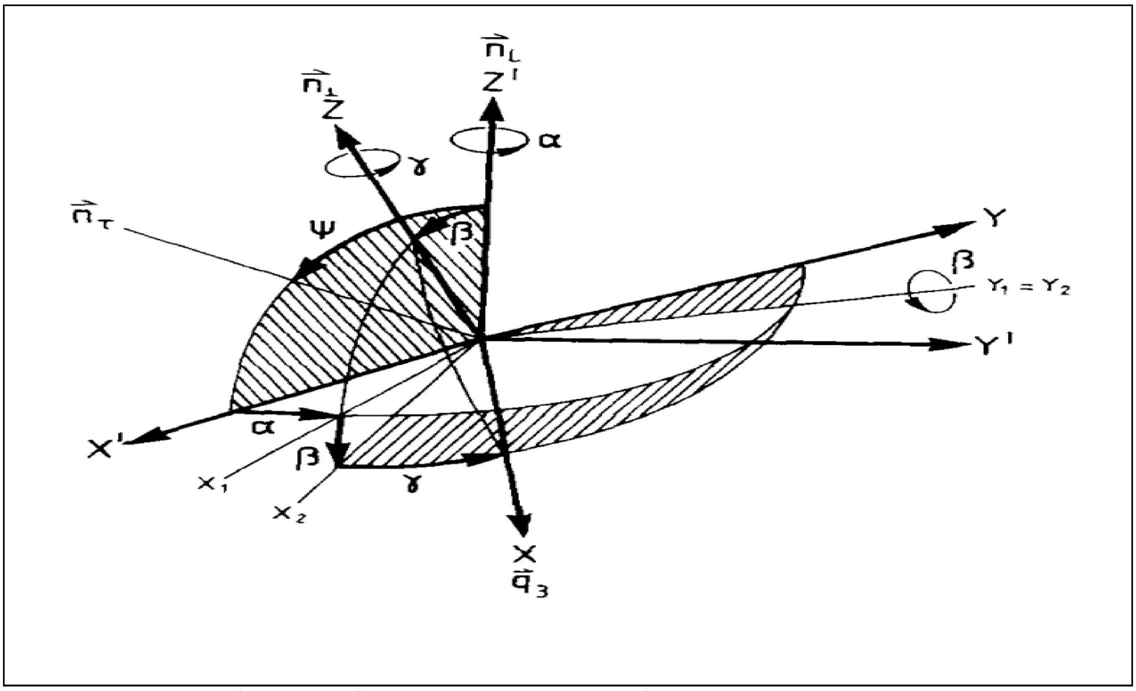

Figure 4.4: Definition of the Euler angles α, β and γ relating the two coordinates S and S' [18].

frame). The x' axis of the S' will be chosen such that the direction of flight of the τ as seen from the hadronic rest frame (n_τ) is in the (x', z')-plane with $n_L \times n_\tau / |n_L \times n_\tau|$ pointing along the Oy' direction[1].

The S frame is chosen to allow for a simple description of the hadron tensor. The (x,y)-plane is aligned with the hadron momentum, with normal to the hadronic plane given by $n_\perp = q_1 \times q_2 / |q_1 \times q_2|$ pointing along the Oz. The Ox axis is defined by the direction of the $\hat{q_3} = \vec{q_3}/|\vec{q_3}|$.

The two frames are related by a Euler rotation R from Fig. 1 given as $x = R(\alpha, \beta, \gamma)x'$ with

[1] One thing to note is that the direction of the spin vector(s_τ) of τ also lies in the (x', z')-plane if the τ polarization vector s is aligned with the τ direction of flight in the laboratory frame

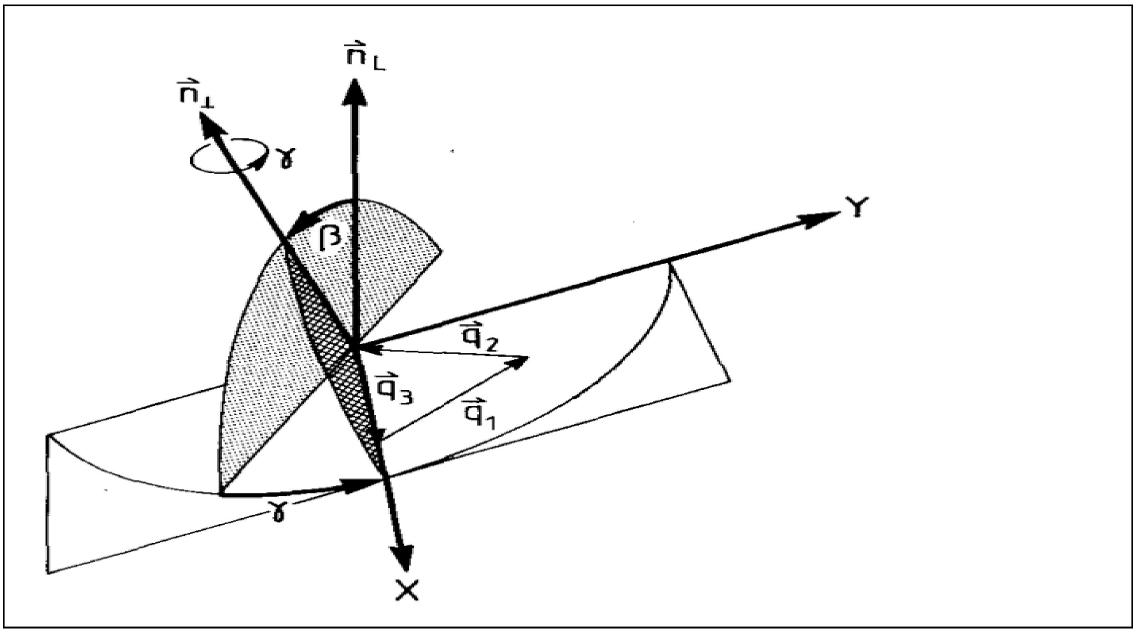

Figure 4.5: Definition of the polar angle β and azimuthal angle γ. Here β is the angle between $n_\perp$ and n_L and γ is the angle between the $(n_L, n_\perp)$ plane and the $(n_L, \hat{q}_3)$ plane [18].

$$R(\alpha, \beta, \gamma) =$$

$$\begin{pmatrix} \cos\alpha\cos\beta\cos\gamma - \sin\alpha\sin\gamma & \sin\alpha\cos\beta\cos\gamma + \cos\alpha\sin\gamma & -\sin\beta\cos\gamma \\ -\cos\alpha\cos\beta\sin\gamma - \sin\alpha\cos\gamma & -\sin\alpha\cos\beta\sin\gamma + \cos\alpha\cos\gamma & \sin\beta\sin\gamma \\ \sin\beta\cos\alpha & \sin\beta\sin\alpha & \cos\beta \end{pmatrix} \tag{4.46}$$

The azimuthal angle α is defined as the angle between the two planes (n_L, n_τ) and $(n_L, n_\perp)$. The angle β is defined as the angle between n_L and $n_\perp$ and the Euler angle γ corresponds to a rotation around the $n_\perp$ and determines the orientation of the hadrons with their production plane, i.e., the angle between $O(z', z)$ plane ($= (n_L, n_\perp)$ plane) and the O(z,x) plane ($= (n_\perp, q_3)$ plane). A more informative definition of the angles γ and β is repeated in Fig. 4.5 below.

As proposed in Reference [18], the angles β and γ are observable in the reaction $e^+e^- \to \tau^+\tau^- \to \nu 3h$ even if we cannot reconstruct the τ rest frame and the neutrino escapes. For $0 \le \alpha \le 2\pi$, $0 \le \beta \le \pi$ and $0 \le \gamma \le 2\pi$ we have

$$\cos\beta = n_L \cdot n_\perp \ , \tag{4.47}$$

$$\cos\gamma = -\frac{n_L \cdot \hat{q}_3}{|n_L \times n_\perp|} \ , \tag{4.48}$$

$$\sin\gamma = \frac{(n_L \times n_\perp) \cdot \hat{q}_3}{|n_L \times n_\perp|} \ , \tag{4.49}$$

$$\cos\alpha = \frac{(n_L \times n_\tau) \cdot (n_L \times n_\perp)}{|n_L \times n_\tau||n_L \times n_\perp|} \ , \tag{4.50}$$

$$\sin\alpha = -\frac{n_\tau \cdot (n_L \times n_\perp)}{|n_L \times n_\tau||n_L \times n_\perp|} \ . \tag{4.51}$$

From Eq. (5.44) we can expand the lepton and hadron tensors as

$$L_{\mu\nu}H^{\mu\nu} = \sum_X L_X W_X = 2(m_\tau^2 - Q^2)\sum_X \bar{L}_X W_X \ , \tag{4.52}$$

where $X = \{A, B, ...I, SA, SB, ..., SG\}$ where the terms relevant for the CP violations W_F, W_G, W_H, W_I, W_{SB}, W_{SC}, W_{SD}, W_{SE}, W_{SF} and W_{SG} are given as [18],

$$W_F = 2x_4[x_1 Im(F_1F_3^*) + x_2 Im(F_2F_3^*)] \ , \tag{4.53}$$

$$W_G = -2x_4[x_1 Re(F_1F_3^*) + x_2 Re(F_2F_3^*)] \ , \tag{4.54}$$

$$W_H = 2x_3x_4[Im(F_1F_3^*) - Im(F_2F_3^*)] \ , \tag{4.55}$$

$$W_I = -2x_3x_4[Re(F_1F_3^*) - Re(F_2F_3^*)] \ , \tag{4.56}$$

$$W_{SB} = 2\sqrt{s}[x_1 Re(F_1F_4^*) + x_2 Re(F_2F_4^*)] \ , \tag{4.57}$$

$$W_{SC} = -2\sqrt{s}[x_1 Im(F_1F_4^*) + x_2 Im(F_2F_4^*)] \ , \tag{4.58}$$

$$W_{SD} = 2\sqrt{s}x_3[Re(F_1F_4^*) - Re(F_2F_4^*)] \ , \tag{4.59}$$

$$W_{SE} = -2\sqrt{s}x_3[Im(F_1F_4^*) - Im(F_2F_4^*)] \ , \tag{4.60}$$

$$W_{SF} = -2\sqrt{s}x_4 Im(F_3F_4^*) \ , \tag{4.61}$$

$$W_{SG} = -2\sqrt{s}x_4 Re(F_3F_4^*) \ , \tag{4.62}$$

where { W_F, W_G, W_H, W_I } are the terms which can probe the CP violation coming from NP of vector type and { W_{SB}, W_{SC}, W_{SD}, W_{SE}, W_{SF}, W_{SG} } are the terms which can probe CP violation coming from NP of scalar type.

4.3 Experimental Searches For The CP Violation in Hadronic τ Decays

Electric Dipole Moment (EDM)

The present experimental bound on the CP(T) violation related to τ lepton are derived from the measurements of electric dipole moment (EDM) and CP violation in hadronic τ decays. The EDM of the τ lepton is given by the effective operator

$$\frac{i}{2} e d_\tau \bar{\tau} \sigma_{\mu\nu} \gamma_5 \tau F_{\mu\nu} \ , \tag{4.63}$$

where d_τ is the EDM of the τ and $F^{\mu\nu}$ is the electromagnetic tensor. In the non-relativistic limit, using the two component spinor ϕ, the Eqs(4.63) reduces to [14]

$$i e d_\tau \phi^\dagger S \phi \cdot E \ , \tag{4.64}$$

where $S = \frac{\sigma}{2}$ and E is the electric field strength. Non zero value of d_τ implies T violation because the operator is odd under T, i.e.,

$$T^{-1} E T = E \text{ and } T^{-1} \phi^\dagger S \phi T = -\phi^\dagger S \phi \ . \tag{4.65}$$

The latest experimental bounds on the electric dipole moment (edm) from PDG averaged data gives,

$$\begin{aligned} Re(e\ d_\tau) &< 0.50 \times 10^{-17} (e\ cm) \ , \\ Im(e\ d_\tau) &< 1.1 \times 10^{-17} (e\ cm) \ . \end{aligned} \tag{4.66}$$

CLEO Searches of CP Asymmetries in τ Decays

CLEO collaboration have searched for CP violation in $\tau \to \pi\pi\nu_\tau$ [17] and $\tau \to K_S \pi \nu_\tau$ [10], but as pointed out by the authors in Reference [15], the CP violation in the $\tau \to \pi\pi\nu_\tau$ is very small so it is not expected to observe CP violation in the $\tau \to \pi\pi\nu_\tau$ decay. Now for the CP violation in $\tau \to K_S \pi \nu_\tau$ coming from New Physics (NP) of scalar type, CLEO detector at CESR operating near Υ(4S) resonance collected data corresponding

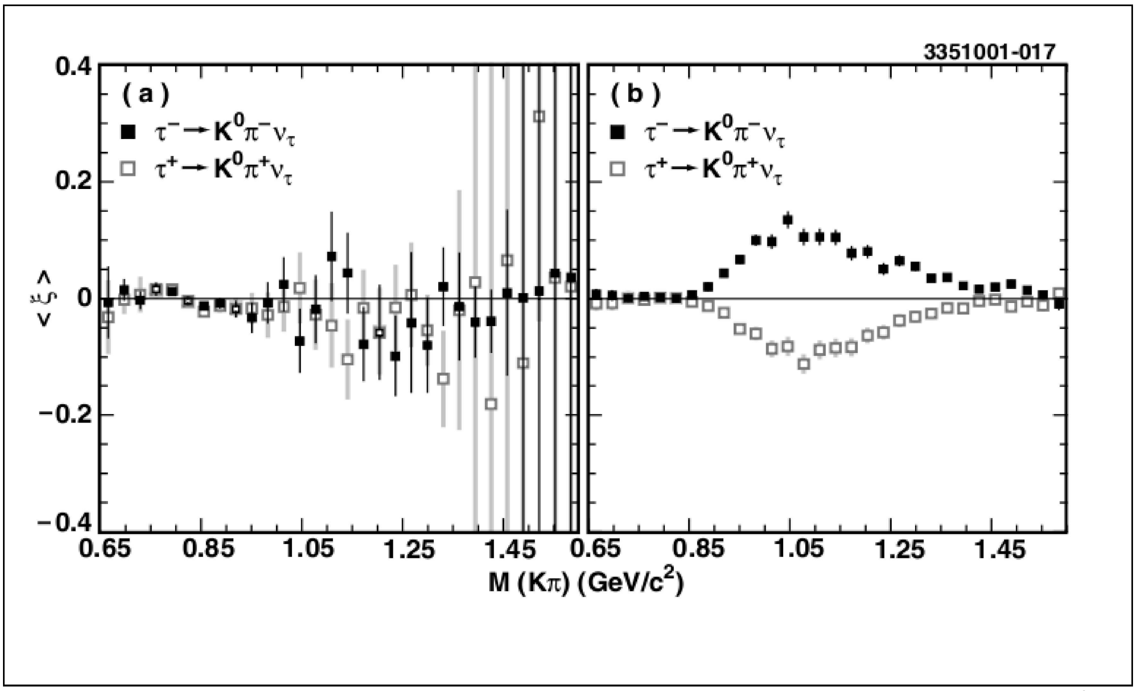

Figure 4.6: Average value of the optimal observable as a function of the $(K_S^0\pi)$ invariant mass for the data (left) and Monte Carlo (right) with the maximum CP violation i.e., $\text{Im}(\Lambda) = 1$ where Λ is the complex coupling of the new scalar [10].

to a total integrated luminosity of about 13.3 fb^{-1} which contains about 12.2 million $\tau^+\tau^-$ pairs. The backgrounds are estimated by analyzing samples of the Monte Carlo (MC) samples. The generation of τ pair production and decay is modeled by KORALB event generator, modified to include the charged scalar contribution to the $\tau \to K_S\pi\nu_\tau$ decay. The detector response is simulated with a GEANT-based Monte Carlo (MC). The main observable that they have used is the "optimal" observable $\xi = P_{odd}/P_{even}$ as defined in Eq. 4.2.

In Figure 4.6 the plot of $\langle\xi\rangle$ is shown separately for τ^+ and τ^- as a function of $(K_S^0\pi)$ invariant mass for the data (left plot) and for the Monte Carlo (MC) with maximum CP violation (right plot). A difference between the $\langle\xi\rangle$ distributions of τ^+ and τ^- would signal a CP violating New Physics (NP). However CLEO data shows no difference in the $\langle\xi\rangle$ distributions of τ^+ and τ^-, therefore, no CP violation is observed. For a restricted range of $(K_S^0\pi)$ mass between (0.85 and 1.45 GeV), the limits on the imaginary part of the scalar coupling are given as

$$Im(\Lambda) = (-0.046 \pm 0.044 \pm 0.019)(1 \pm 0.15) \ , \tag{4.67}$$

where the first error is statistical and the second error is additive systematic and the multiplying term is the multiplicative systematic error. The corresponding limits are given as

$$-0.172 \ < \ Im(\Lambda) \ < \ 0.067 \tag{4.68}$$

at 90% C.L.

Belle Searches of CP Asymmetries in τ Decays

Belle also searched for CP violation coming from New Physics (NP) of scalar type [3]. Using the Belle detector at KEKB asymmetric-energy e^+e^- collider operating at $\Upsilon(3S)$, $\Upsilon(4S)$ and $\Upsilon(5S)$ resonances 699 fb^{-1} of data were collected. The signal and backgrounds from $\tau^+\tau^-$ are simulated by KKMC/TAULA [3] and detector response is simulated by a GEANT3 based program. The CP observable they have used to search for CP violation is [3, 7],

$$\begin{aligned} A^i_{CP} &= \frac{\int\int\int \cos\beta\cos\psi(\frac{d\Gamma_{\tau^+}}{d\omega_i} - \frac{d\Gamma_{\tau^-}}{d\omega_i})d\omega_i}{\frac{1}{2}\int\int\int(\frac{d\Gamma_{\tau^+}}{d\omega_i} + \frac{d\Gamma_{\tau^-}}{d\omega_i})d\omega_i} \\ &\simeq \langle\cos\beta\cos\psi\rangle^i_{\tau^+} - \langle\cos\beta\cos\psi\rangle^i_{\tau^-} \ , \end{aligned} \tag{4.69}$$

where $d\omega_i = dQ_i^2 d\cos\theta d\cos\psi$ in bin i. The angle β is the angle between the direction of the K_S^0 and the direction of the e^+e^- center of the mass (c.m) frame, both as measured in the hadronic rest frame. The angle θ is the angle between direction opposite to the direction of the center of the mass (c.m) frame and the direction of the hadronic rest frame as seen from the τ rest frame and given as in Eq. 4.21. The angle ψ denote the angle between the direction of the center of the mass (c.m) frame and the direction of the τ as seen from the hadronic rest frame and given by Eq. 4.22. The background subtracted asymmetry is shown in Fig. 4.7 with the statistical and systematic errors added in quadrature where the vertical error bars are the systematic errors added in quadrature [3].

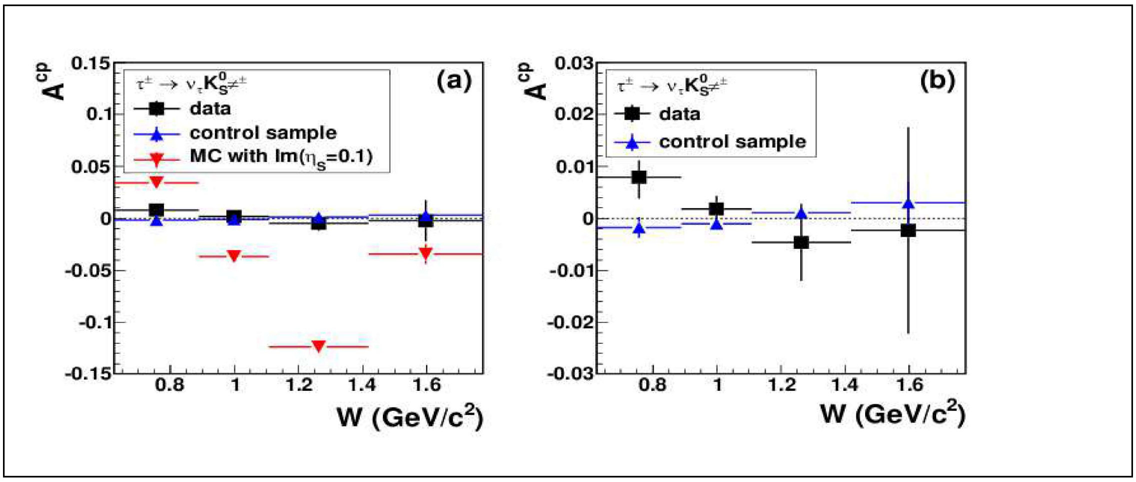

Figure 4.7: (a) The measured CP violation asymmetry after background subtraction are made (squares). (b) The expanded view where the vertical scale is reduced by a factor of 5 [3].

In Fig. 4.7 the CP asymmetry measured in the controlled sample is indicated by the blue triangles (statistical errors only) and the inverted red triangles show the expected asymmetry for the $Im(\eta_S) = 0.1$ $[Re(\eta_S) = 0]$ where η_S is the complex scalar coupling from the NP. Comparing with CLEO's notation in Eq. 4.3 we have $\eta_S = 1.1\Lambda$ [3]. Data from the Belle search for CP violation in the $\tau^{\pm} \to K_S^0 \pi^{\pm} \nu_\tau$ decays by analyzing asymmetries in the angular distributions of τ^+ with respect to τ^- show no significant CP asymmetry. Therefore Belle have set an upper limit for the CP violation parameter $Im(\eta_S)$ in the range of $|Im(\eta_S)| < 0.026$ or better, depending on the parametrization used to describe the hadronic form factors, at 90% confidence level.

Chapter 5

Observed Sign Anomaly in $A_{cp}(\tau \to K_S\pi\nu_\tau)$

5.1 Introduction

Apart from the CP phases that may arise in the neutrino mixing matrix, the decays of the tau lepton may allow us to explore nonstandard CP-violating interactions. Various experimental groups have been involved in exploring CP violation in tau decays in the last decade or more. In 2002, the CLEO collaboration [2], and more recently the Belle Collaboration[3], studied the angular distribution of the decay products in $\tau \to K_s\pi\nu_\tau$ in search of CP violation from model dependent new scalar particle exchanges; however, neither study observed any CP asymmetry. The BaBar collaboration for the first time reported a sign anomaly in the measurement of CP asymmetry in the $\tau \to K_S\pi\nu_\tau$ decay [20]:

$$A_\tau = \frac{\Gamma(\tau^+ \to K_S\pi^+\bar{\nu_\tau}) - \Gamma(\tau^- \to K_S\pi^-\nu_\tau)}{\Gamma(\tau^+ \to K_S\pi^+\bar{\nu_\tau}) + \Gamma(\tau^- \to K_S\pi^-\nu_\tau)} = -(0.36 \pm 0.23 \pm 0.11)\% \ . \tag{5.1}$$

The (BaBar, BELLE, CLEO and FOCUS) [37] collaborations have reported the measured CP asymmetry in the $D \to K_S\pi$ decay as:

$$A_D = \frac{\Gamma(D^+ \to K_S\pi^+) - \Gamma(D^- \to K_S\pi^-)}{\Gamma(D^+ \to K_S\pi^+) + \Gamma(D^- \to K_S\pi^-)} = -(0.54 \pm 0.14)\% \ , \tag{5.2}$$

where the value given is the average of the four measurements. Now if only CP asymmetry in these decay modes is the SM one coming form $K^0 - \bar{K}^0$ mixing then as argued by Grossman and Nir in Reference [37], the sign of the CP asymmetry observed in the A_τ and A_D should be opposite. In Section 5.2 we will reproduce the Grossman and Nir's basic argument of why the sign in the A_τ and A_D should be opposite and especially why the sign in the A_τ should be positive, contrary to BaBar's observation of negative sign.

5.2 CP Violation From $K^0 - \bar{K}^0$ Mixing in SM

In SM there is no CP violation in the leptonic sector and the only source of CP violation in the quark sector is the one phase in the CKM mixing matrix. So the measured CP asymmetry in $\tau^\pm \to K_S\pi^\pm\nu_\tau$ and $D^\pm \to K_S\pi^\pm$ from the B-factories[BaBar,BELLE,CLEO and FOCUS] comes mainly from the CP violation in the $K^0 - \bar{K}^0$ mixing in SM. As Grossman and Nir pointed out [37], BaBar's measurement of CP asymmetry in the $\tau^\pm \to K_S\pi^\pm\nu_\tau$:

$$A_\tau = -(0.36 \pm 0.23 \pm 0.11)\% \tag{5.3}$$

and (BaBar, BELLE, CLEO and FOCUS)[37] collaborations measurements of CP violation in the $D^\pm \to K_S\pi^\pm$ as:

$$A_D = -(0.54 \pm 0.14)\% \tag{5.4}$$

are in conflict with the SM prediction of CP violation in these decays coming from the $K^0-\bar{K}^0$ mixing alone, as $K^0-\bar{K}^0$ mixing predict these two decay modes to have opposite sign. The reason why the $\tau^\pm \to K_S\pi^\pm\nu_\tau$ and $D^\pm \to K_S\pi^\pm$ should have opposite sign is that the $\tau^+(\tau^-)$ decay initially to $K^0(\bar{K}^0)$ while $D^+(D^-)$ decay initially to $\bar{K}^0(K^0)$ state neglecting the color and doubly Cabibbo suppressed decay $D^+ \to K^0\pi^+$. Since the intermediate K_S state is not directly observed in the experiments but it is reconstructed via the $\pi^+\pi^-$ final state with $m_{\pi\pi} \approx m_{K_S}$ and time difference between τ and D decay and K decay time $t \approx \tau_S$ where τ_S is the lifetime of the K_S. Thus, the CP asymmetry from $K^0(\bar{K}^0)$ mixing in SM depends on the integrated decay times given as

$$A_\tau(t_1, t_2) = -A_D(t_1, t_2) = A_\epsilon(t_1, t_2), \tag{5.5}$$

with

$$A_\epsilon(t_1, t_2) = \frac{\int_{t_1}^{t_2}[\Gamma(K^0(t) \to \pi\pi) - \Gamma(\bar{K}^0(t) \to \pi\pi)]}{\int_{t_1}^{t_2}[\Gamma(K^0(t) \to \pi\pi) + \Gamma(\bar{K}^0(t) \to \pi\pi)]} \tag{5.6}$$

where the $K^0(t)$ $(\bar{K}^0(t))$ are time evolved states at time t from initial pure states $K^0(0)$ $(\bar{K}^0(0))$. Although the theoretical prediction depends on the t_1, t_2 and on the details of the experiment, the fact that $A_\tau(t_1, t_2)$ and $A_D(t_1, t_2)$ are predicted to have opposite sign by SM, while experimental measurements [37] of these observables carry the same sign is interesting even though the naive expectation that $A_\tau = -A_D$ is excluded at 3.3σ. It is very difficult to assess the significance of $A_\tau \neq -A_D$ until the dependence of theoretical prediction on the t_1, t_2 and on details of the experiment are properly taken into consideration. In the next section the details of the phenomenological prediction of CP violation from $K^0 - \bar{K}^0$ mixing in SM is derived.

Phenomenology of $K^0 - \bar{K}^0$ Mixing in SM

The states $|K^0\rangle$ and $|\bar{K}^0\rangle$ are the flavor eigenstates and the states $|K_S\rangle$ and $|K_L\rangle$ are the mass eigenstates of masses m_S and m_L and widths Γ_S and Γ_L respectively. The states $|K_S\rangle$ and $|K_L\rangle$ are linear combinations of the flavor eigenstates $|K^0\rangle$ and $|\bar{K}^0\rangle$ are given as

$$|K_{S,L}\rangle = p|K^0\rangle \pm q|\bar{K}^0\rangle \ . \tag{5.7}$$

where the p and q are such that in the case of no CP violation in the mixing we have $p = q = \frac{1}{\sqrt{2}}$. In the analysis of CP violation from the $K^0 - \bar{K}^0$ mixing the useful variables turn out to be the average and difference in mass and width given as
$m = \frac{m_S + m_L}{2}, \ \Gamma = \frac{\Gamma_S + \Gamma_L}{2}, \ \Delta m = m_L - m_S,$

$$\Delta\Gamma = \Gamma_L - \Gamma_S, \ x = \frac{\Delta m}{\Gamma}, \ y = \frac{\Delta\Gamma}{2\Gamma} \ . \tag{5.8}$$

Now if CP is conserved then only $K_S \to \pi\pi$ is allowed, but in presence of CP violating dynamics in the $K^0 - \bar{K}^0$ mixing both $K_S \to \pi\pi$ and $K_L \to \pi\pi$ are allowed. Then the decay amplitudes into a final state $\pi\pi$ are defined as

$$A_{S,L} = \langle \pi\pi | H | K_{S,L} \rangle \ . \tag{5.9}$$

The relevent CP violating parameters are defined as

$$\frac{|p|^2 - |q|^2}{|p|^2 + |q|^2} \approx 2Re(\epsilon) \text{ with } \tfrac{A_L}{A_S} \approx \epsilon \ , \tag{5.10}$$

where terms of order $|\epsilon|^2$ and $\frac{\epsilon'}{\epsilon}$ are ignored. The ϵ' being a measure of direct CP violation. Now for the difference and sum of the time dependent decay rates of $K^0(t) \to \pi\pi$ and $\bar{K}^0(t) \to \pi\pi$ are given as [37]

$$\frac{D_{\pi\pi}(t)}{N} = -2Re(\epsilon)(e^{-\Gamma_S t} + |\epsilon|^2 e^{-\Gamma_L t}) + 2e^{-\Gamma t}(Re(\epsilon)\cos(\Delta m t) + Im(\epsilon)\sin(\Delta m t)) \ , \tag{5.11}$$

and

$$\frac{S_{\pi\pi}(t)}{N} = e^{-\Gamma_S t} + |\epsilon|^2 e^{-\Gamma_L t} - 4Re(\epsilon)e^{-\Gamma t}(Re(\epsilon)\cos(\Delta m t) + Im(\epsilon)\sin(\Delta m t)) \ , \tag{5.12}$$

where $N = |A_S|^2 \frac{|p|^2+|q|^2}{4|p|^2|q|^2}$. The plots of ratio between the second term (interference term) and the first term (non interference term) in $D_{\pi\pi}(t)$ given as [37]

$$R(t) = -\frac{e^{-\Gamma t}(\cos(\Delta m t) + \frac{Im(\epsilon)}{Re(\epsilon)}\sin(\Delta m t))}{e^{-\Gamma_S t} + |\epsilon|^2 e^{-\Gamma_L t}} \ , \tag{5.13}$$

are shown in Fig. 5.1 below.

From Fig. 5.1 the key observations made in [37] are :

1. The interference term is not negligible even at the very early times, e.g at $t = 0$, $R = -1$.

2. In the approximation $Re(\epsilon) \approx Im(\epsilon)$ and $x \approx -y$, the ratio changes sign when $\tan(\frac{1}{2}t/\tau_S) = -1$ that is $t/\tau_S = 3\pi/2 + 2n\pi$ for n = (0,1,2.....).

3. For times early enough that pure K_L terms can be negleted ($t \ll 12\tau_S$), R reaches a minimum at $t/\tau_S = \pi$, $R = -e^{\pi/2}$, and maximum at $t/\tau_S = 3\pi$, $R = +e^{3\pi/2}$.

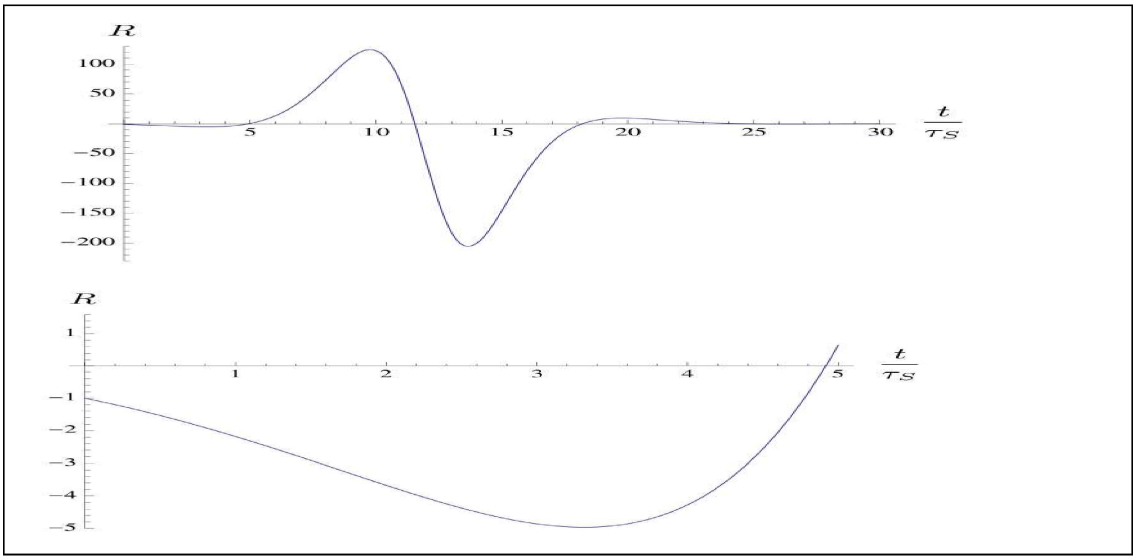

Figure 5.1: Plot of R(t) in Eq. (5.13) as a function of time in units of τ_S (top) and the zoomed into the short region of time interval of 0 to 5 in unit of τ_S (bottom) taken from Reference [37].

The experimentally measured CP asymmetry is thus given by the convolution of the bare asymmetry with a time dependent funtion F :

$$A_\epsilon(t) = \frac{\int_0^\infty F(t)D_{\pi\pi}(t)dt}{\int_0^\infty F(t)S_{\pi\pi}(t)dt} , \tag{5.14}$$

where the convolution function $F(t)$ parametrize the dependence of the experimentally measured CP asymmetry on the kaon decay time. The final measurement is sensitive to the experimental cut i.e., it depends on the K_S selection procedure employed. It is reasonable to approximate the function F(t) given as a step function, $F(t) = 1$ for $t_1 \ll t \ll t_2$ else 0. For the case where $t_2 \ll \tau_L$, we have [37]

$$A_\epsilon(t_2, t_1) = -2Re(\epsilon)\left\{1 - \frac{\int_{t_1}^{t_2} dte^{-\Gamma t}(\cos(\Delta m t) + \frac{Im(\epsilon)}{Re(\epsilon)}\sin(\Delta m t))}{\int_{t_1}^{t_2} dte^{-\Gamma_S t}}\right\} , \tag{5.15}$$

where terms of $\mathcal{O}(\epsilon^2)$ are neglected. In the case where direct CP violation can be neglected, by using the model independent relation [37]

$$\frac{Im(\epsilon)}{Re(\epsilon)} = -\frac{x}{y} , \tag{5.16}$$

they obtained

$$A_\epsilon(t_2,t_1) = -2Re(\epsilon)\left\{1 - \frac{2(1-x^2/y)}{1+x^2}\frac{e^{-\Gamma t_1}\cos\Delta m t_1 - e^{-\Gamma t_2}\cos\Delta m t_2}{e^{-\Gamma_S t_1} - e^{-\Gamma_S t_2}} + \frac{2(x-x/y)}{1+x^2}\frac{e^{-\Gamma t_1}\sin\Delta m t_1 - e^{-\Gamma t_2}\sin\Delta m t_2}{e^{-\Gamma_S t_1} - e^{-\Gamma_S t_2}}\right\}. \tag{5.17}$$

In the region where we can take the approximation as $t_1 \ll \tau_S$ and $\tau_S \ll t_2 \ll \tau_L$ so that we can take $e^{-\Gamma_S t_1} = 1$, $e^{-\Gamma_S t_2} = 0$ and $\cos(\Delta m t_1) = 1$ and also using $y \approx -1$ they obtained

$$A_\epsilon(t_1 << \tau_S, \tau_S << t_2 << \tau_L) = +2Re(\epsilon) \approx 3.3\times 10^{-3}, \tag{5.18}$$

where the numerical value is the experimental value. In Figs. 5.2 and 5.3 the dependence of $A_\epsilon(t_1 t_2)$ on the choice of the t_1 and t_2 is shown. In Fig. 5.2 a plot of the dependence of $A_\epsilon(t_1 t_2)/(2Re(\epsilon))$ as a function of t_1 for $t_2 = 10\tau_S$ and in Fig. 5.3 a plot of the dependence of the $A_\epsilon(t_1 t_2)/(2Re(\epsilon))$ as a function of t_2 for $t_2 = \tau_S/10$ are shown as in Reference [37]. CP asymmetry of $\mathcal{O}(10^{-3})$ are predicted in the $\tau^\pm \to K_S\pi^\pm\nu_\tau$ and $D^\pm \to K_S\pi^\pm$ in the SM as a result of the CP violation in the $K^0 - \bar{K}^0$ mixing. Any deviation from this prediction would imply direct CP violation in τ and/or D decays.

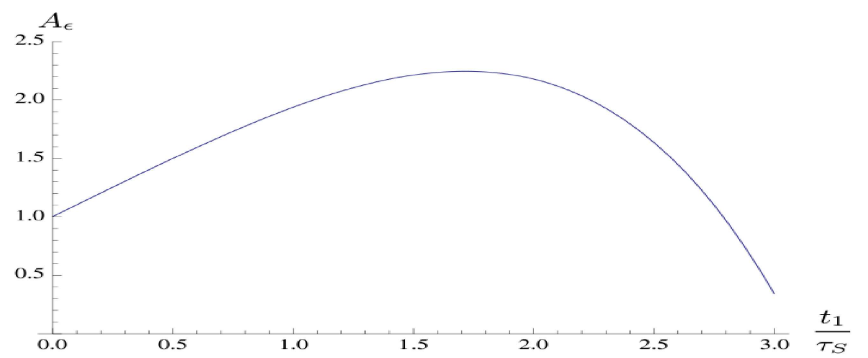

Figure 5.2: Plot of $A_\epsilon(t_1,t_2)$ given in Eqs(5.17) as a function of t_1/τ_S in units of $[2Re(\epsilon)]$ for $t_2 = 10\tau_S$ taken from the Reference [37].

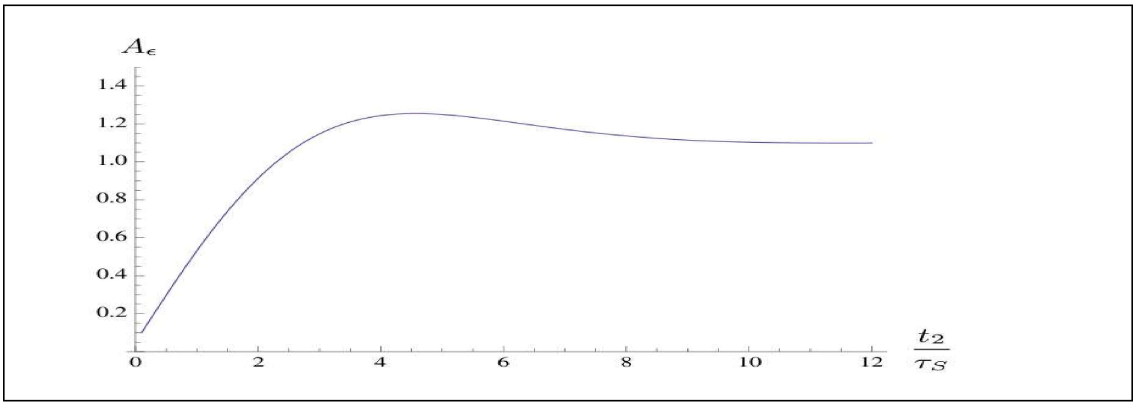

Figure 5.3: Plot of $A_\epsilon(t_1, t_2)$ given in Eqs(5.17) as a function of t_2/τ_S in units of $[2Re(\epsilon)]$ for $t_1 = \tau_S/10$ taken from the Reference [37].

5.3 $A_{cp}(\tau \to K_S\pi\nu_\tau)$ in presence of new tensor Interaction

The BABAR collaboration's observation [20] of the sign anomaly in the integrated decay rate asymmetry $A_{cp}(\tau \to K_S\pi\nu_\tau)$ of

$$A_{cp}^{Exp} = (-0.36 \pm 0.23 \pm 0.11)\% \ , \tag{5.19}$$

different than expected from the SM as argued in Section 5.2 may be a sign of a new source of CP violation other than the SM one coming from the $K^0 - \bar{K}^0$ mixing in hadronic τ decays. If the result stand in the more accurate future experimentations, it is very interesting because it would be the first time a CP violation in the leptonic sector was ever discovered outside of the possible CP violation in the neutrino mixing. Naively one may expect that the simplest way to account for the observed anomaly would be to introduce a direct CP violation via a new CP violating charged scalar exchange. However, it turns out that the charged scalar type of exchange may contribute in the angular distributions, but its mixing with SM term in the integrated decay rate goes to zero. Now the next candidate of NP would be a new CP violating charged vector exchange, but CP violation from vector type NP will be observable only if both vector current and axial vector currents contribute to the same final states [6, 7]. Since in two pseudo scalar meson final states only vector current can contribute due to parity conservation of strong interaction, vector type of NP can contribute in general to CP violation in three or more pseudo scalar meson final states but not in two pseudo scalar meson final states such as $K_s\pi$. Now the only possibility left is tensor type of NP.

Effective Hamiltonian And Decay Rates

With the assumption that all neutrinos are left handed, we propose the most general effective Hamiltonian containing all possible four fermion interaction operators that can contribute to $\tau \to K_S\pi\nu_\tau$ as given by:

$$H_{eff} = \frac{4G_F}{\sqrt{2}}V_{us}[(\delta_{l3}+C^{\tau}_{V_1})\mathcal{O}^{\tau}_{V_1}+C^{\tau}_{V_2}\mathcal{O}^{l}_{V_2}+C^{\tau}_{S_1}\mathcal{O}^{\tau}_{S_1}+C^{\tau}_{S_2}\mathcal{O}^{\tau}_{S_2}+C^{\tau}_{T}\mathcal{O}^{\tau}_{T}]+h.c.\ , \quad (5.20)$$

with the operators given by

$$\mathcal{O}^{\tau}_{V_1} = (\bar{s_L}\gamma^{\mu}u_L)(\bar{\nu_L}\gamma_{\mu}\tau_L)\ , \quad (5.21)$$

$$\mathcal{O}^{\tau}_{V_2} = (\bar{s_R}\gamma^{\mu}u_R)(\bar{\nu_L}\gamma_{\mu}\tau_L)\ , \quad (5.22)$$

$$\mathcal{O}^{\tau}_{S_1} = (\bar{s_R}u_L)(\bar{\nu_L}\tau_R)\ , \quad (5.23)$$

$$\mathcal{O}^{\tau}_{S_2} = (\bar{s_L}u_R)(\bar{\nu_L}\tau_R)\ , \quad (5.24)$$

$$\mathcal{O}^{\tau}_{T} = (\bar{s_L}\sigma^{\mu\nu}u_R)(\bar{\nu_L}\sigma_{\mu\nu}\tau_R)\ . \quad (5.25)$$

Since we are concerned with CP violation in $\tau \to K_S\pi\nu_\tau$, we can set the $C^{\tau}_{V_1}$ and $C^{\tau}_{V_2}$ equal to zero for simplicity as these coefficients will not contribute in CP violation in two meson final states as argued earlier. As we mentioned earlier and argued in a paper of ours [8] that in the integrated decay rate asymmetry the contribution from the charged scalars goes to zero, so the only terms left are the SM term and the tensor term.

Decay Rate of $\tau \to K_S \pi \nu_\tau$ in SM

In the SM the $\tau \to K_S \pi \nu_\tau$ decay rate can be expressed as:

$$d\Gamma_{SM}(\tau \to K\pi\nu) = \frac{1}{2m_\tau}\frac{G_F^2}{2}V_{us}^2 \mathcal{L}_{\mu\nu}\mathcal{H}^{\mu\nu} dPS^{(3)} \tag{5.26}$$

where

$$\mathcal{L}_{\mu\nu} = [\bar{\nu}_\tau \gamma_\mu (1-\gamma_5)\tau]\,[\bar{\nu}_\tau \gamma_\nu (1-\gamma_5)\tau]^\dagger \tag{5.27}$$

and

$$\mathcal{H}^{\mu\nu} = \mathcal{J}^\mu (\mathcal{J}^\nu)^\dagger \tag{5.28}$$

where

$$\mathcal{J}^\mu = \langle K(q_1)\pi(q_2)|V^\mu(0)|0\rangle. \tag{5.29}$$

The hadronic current can be parametrized in terms of the vector and scalar form factors as:

$$\mathcal{J}^\mu = F_V^{K\pi}(Q^2)\big(g^{\mu\nu} - \frac{Q^\mu Q\nu}{Q^2}\big)(q_1 - q_2)_\nu + \frac{(m_K^2 - m_\pi^2)}{s} F_S^{K\pi} Q^\nu \tag{5.30}$$

where $Q^\mu = (q_1 + q_2)^\mu$ and in the hadronic rest frame the decay rate can be expressed as:

$$\begin{aligned}\frac{d\Gamma_{SM}(K\pi)}{ds} = \frac{G_F^2 V_{us}^2 m_\tau^3}{3 \times 64\pi^3}\frac{1}{s^{\frac{3}{2}}}\left(1 - \frac{s}{m_\tau^2}\right)^2 \left(1 + \frac{2s}{m_\tau^2}\right) \times P(s) \\ \left\{P(s)^2 |F_V|^2 + \frac{3(m_k^2 - m_\pi^2)^2}{4s(1+\frac{2s}{m_\tau^2})}|F_S|^2\right\}\end{aligned} \tag{5.31}$$

where

$$P(s) = |\vec{q_1}| = \frac{1}{2\sqrt{s}}\sqrt{[s - (m_k + m_\pi)^2]\,[s - (m_k - m_\pi)^2]} \tag{5.32}$$

is the momentum of the K in the $K\pi$ rest frame and s is the $K\pi$ invariant mass squared i.e $s = Q^2$. The vector form factor can be parameterized by $K^*(892)$,$K^*(1410)$ and $K^*(1680)$ meson amplitudes as given in Reference [9]:

$$F_V = \frac{1}{1+\beta+\chi}\left[BW_{K^*(892)}(s) + \beta BW_{K^*(1410)}(s) + \chi BW_{K^*(1680)}(s)\right] \tag{5.33}$$

where β and χ are the complex coefficients for the fractions of $K^*(1410)$ and $K^*(1680)$ resonances respectively and $BW_R(s)$ is a relativistic Breit-Wigner function for R = $K^*(892)$,$K^*(1410)$ and $K^*(1680)$ given as:

$$BW_R(s) = \frac{M_R^2}{s - M_R^2 + i\sqrt{s}\Gamma_R(s)} \tag{5.34}$$

and

$$\Gamma_R(s) = \Gamma_{0R}\frac{M_R^2}{s}\left(\frac{P(s)}{P(M_R^2)}\right)^{(2l+1)} . \tag{5.35}$$

Here $\Gamma_R(s)$ is the s dependent total width of the resonance and $\Gamma_{0R}(s)$ is the resonance width at its peak and $l = 1$ for the vector states and $l = 0$ for the s-wave part. Similarly the scalar form factor F_S has $K_0^*(800)$ and $K_0^*(1430)$ contributions and is given as:

$$F_S = \kappa\frac{s}{M^2_{K_0^*(800)}}BW_{K_0^*(800)}(s) + \gamma\frac{s}{M^2_{K_0^*(1430)}}BW_{K_0^*(1430)}(s) \tag{5.36}$$

where κ and γ are the real constants that describe the fractional contributions from $K_0^*(800)$ and $K_0^*(1430)$ respectively. As reported by Belle[9], $K^*_{(892)}$ alone is not enough to describe the $K_s\pi$ mass spectrum. It is best explained by $K^*(892)$+$K^*(1410)$+$K^*(800)$ or $K^*(892) + K^*(1430) + K^*(800)$.

Tensorial Term

We now include the contribution from the tensorial operator as it has been already pointed out earlier that scalar and the vectorial operators would not contribute to the integrated decay rate asymmetry and CPV. The key requirement in the relevant context of explaining the observed CPV in integrated $\tau \to K_s\pi\nu_\tau$ decay rate by the tensorial operator is that its coefficient C_T^l from Eq. (6.20) should be complex so that interference of the SM with this tensor amplitude gives the required CP phase. From Eqs. (6.20) & (6.25) the effective Hamiltonian is given as

$$\mathcal{H}_{eff}^T = \frac{4G_F}{\sqrt{2}}V_{us}C_T^\tau(\bar{s_L}\sigma^{\mu\nu}u_R)(\bar{\nu_L}\sigma_{\mu\nu}\tau_R) \tag{5.37}$$

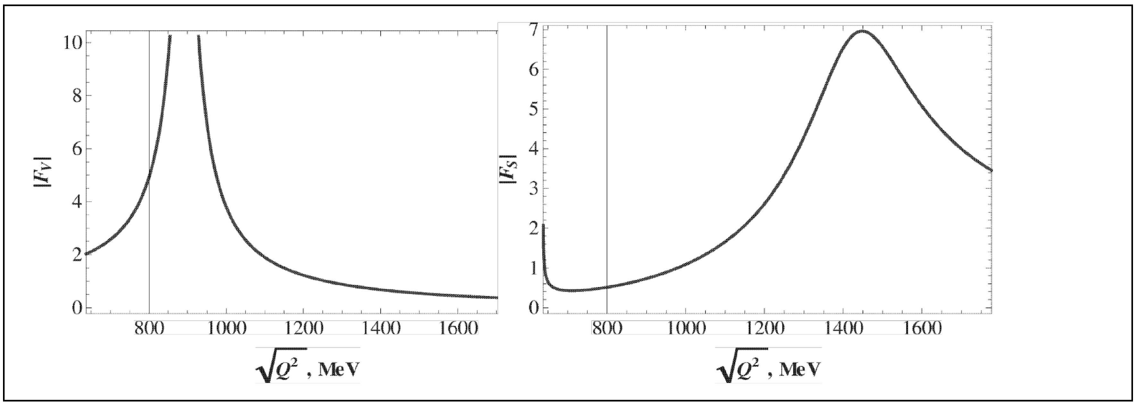

Figure 5.4: Plot of $|F_v|$ as a function of $\sqrt{Q^2}$ (in MeV) with contributions coming from $K^*(892)$ and $K^*(1430)$ (left) and plot of $|F_s|$ as a function of $\sqrt{Q^2}$ (in MeV) [8].

where $\sigma^{\mu\nu} = \frac{i}{2}(\gamma^\mu\gamma^\nu - \gamma^\nu\gamma^\mu)$ and the hadronic current can be expressed as

$$\langle K(q_1)\pi(q_2)|\bar{s}\sigma^{\mu\nu}u|0\rangle = i\frac{F_T}{m_K + m_\pi}(q_1^\mu q_2^\nu - q_2^\mu q_1^\nu) \ , \tag{5.38}$$

where F_T is the tensorial form factor and only tensor terms can contribute due to parity conservation of strong interaction and pseudo-tensor term will not contribute.

Including the contribution from the tensor term to the $\tau \to K_S\pi\nu_\tau$ decay rate

When tensorial term is included the total decay rate is given by

$$d\Gamma = \left(\frac{d\Gamma_{SM}}{ds} + \frac{d\Gamma_{MIX}}{ds} + \frac{d\Gamma_T}{ds}\right) ds \ , \tag{5.39}$$

where the $\frac{d\Gamma_{SM}}{ds}$ is given in Eq. (5.31) and the full angular dependence of the other two terms can be expressed as [19]

$$\frac{d\Gamma_{MIX}}{ds\frac{d\cos\beta}{2}\frac{d\alpha}{2\pi}} = -\frac{G_F^2 V_{us}^2 m_\tau^2}{\pi^3(m_k+m_\pi)4s^{\frac{1}{2}}}(1-\frac{s}{m_\tau^2})^2 P^2 \Big\{ -P\times Re(F_V^\dagger F_T C_T^\tau) + Re(F_S^\dagger F_T C_T^\tau)$$
$$\times[\frac{m_k^2-m_\pi^2}{2\sqrt{s}}]\times(\sin\beta\cos\alpha\sin\psi+\cos\beta\cos\psi)\Big\} \ , \tag{5.40}$$

and

$$\frac{d\Gamma_T}{ds\frac{d\cos\beta}{2}\frac{d\alpha}{2\pi}} = \frac{G_F^2 V_{us}^2 m_\tau^3 |F_T|^2 |C_T^\tau|^2}{(m_k+m_\pi)^2\pi^3 8s^{\frac{1}{2}}}(1-\frac{s}{m_\tau^2})^2 P^2 \Big\{ \frac{P}{2} + \frac{3}{2}(s-m_k^2-m_\pi^2)\frac{(m_k^2-m_\pi^2)}{s^{3/2}}$$
$$\times(\sin\beta\cos\alpha\sin\psi+\cos\beta\cos\psi) - (1-\frac{s}{m_\tau^2})\frac{P}{2}$$
$$\times(\sin\beta\cos\alpha\sin\psi+\cos\beta\cos\psi)^2\Big\} \ , \tag{5.41}$$

where P is the momentum of the K in the $K\pi$ rest frame and the angles α, β are same as defined in Fig. 4.1, and ψ is defined as the angle between direction of flight of the lab frame and the direction of flight of τ as seen from the hadronic rest frame, if τ direction in hadron rest frame in known then we can set $\psi \to 0$. In [8] we have shown, also a derivation is given in Appendix A, that the CP violation coming from the $K-\bar{K}$ mixing and the direct CP violation in $A_{cp}(\tau\to K_S\pi\nu_\tau)$ can be seperated as

$$A_{cp}(\tau\to K_S\pi\nu_\tau) = \frac{A_{cp}^K + A_{cp}^\tau}{1+A_{cp}^K A_{cp}^\tau} \ , \tag{5.42}$$

and also we have

$$Br(\tau\to K_S\pi\nu_\tau) = \frac{(\Gamma^{\tau^+}+\Gamma^{\tau^-})}{2}\tau_\tau \ , \tag{5.43}$$

where A_{cp}^K is the CPV coming from the $K-\bar{K}$ mixing and A_{cp}^τ is the direct CP violation coming from NP particle mediated CP violation at lepton and/or quark vertices and τ_τ is the τ life time. Since both A_{cp}^K and A_{cp}^τ are expected to be small, we can savely ignore terms involving the product of the two.

Results

Case(a):

In *Case*(a) we will assume that F_T is a constant independent of Q^2 and using the Breit Wigner forms of the various resonances contributing to the vector and scalar form factors, the $s(K_S\pi)$(hadronic invariant mass) dependent strong phases were determined. Taking the strong phase of the tensor form factor to be vanishing and its magnitude to be constant, the observables $A_{cp}(\tau \to K_S\pi\nu_\tau)$ and branching fraction were used to determine weak phase and the magnitude of the tensor coupling.

Case(a) I:

Here, we use the combination: K*(892) + K*(1430) + K*(800). Figs. 5.4 and 5.5 show the dependence of the Form Factors and strong phases on $\sqrt{Q^2}$.

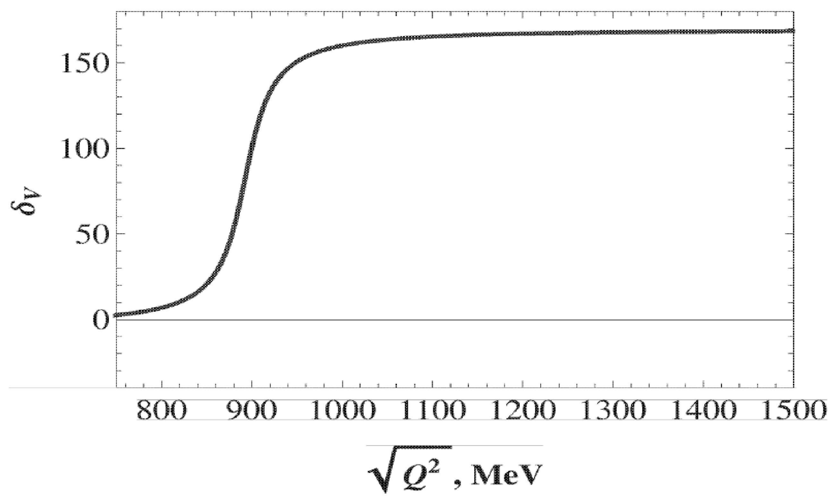

Figure 5.5: Plot of strong phase δ_V = Arg$[F_V]/\pi\times180$ + 180(in degrees) from the combination of $K^*(892)$ and $K^*(1430)$ as a function of $\sqrt{Q^2}$ (in MeV)[8].

Using the average width $\frac{\Gamma^{\tau^+}+\Gamma^{\tau^-}}{2}$, and the CP asymmetry due to direct CP violation $A_{CP} = \frac{\Gamma^{\tau^+}-\Gamma^{\tau^-}}{\Gamma^{\tau^+}+\Gamma^{\tau^-}}$, we can find the solutions for the unknowns $C_T^\tau F_T$ and ϕ, where ϕ is the weak phase from NP. In Table 5.1 we have shown the two feasible solutions. It is pretty obvious that only the first solution is viable, as the NP contribution has to be much smaller than the SM contribution, since there is no glaring evidence of it, other than the unexpected direct CP violation seen. The smaller magnitude of the tensor mod squared and interference term relative to the SM contribution allows the Q^2 distribution of the SM alone to be reasonably consistent with the Belle data.

Sl.No	$C_T^\tau F_T$	$\cos\phi$	$\lvert\frac{A_T}{A_{SM}}\rvert^2$	$\frac{Int(A_T A_{SM})}{A_{SM}^2}$(cos term)	$\frac{Int(A_T A_{SM})}{A_{SM}^2}$(sin term)
(i)	-0.213	-0.816	0.009	0.055	0.039
(ii)	-3.333	0.999	2.234	1.057	0.047

Table 5.1: Table showing the NP contribution parameterized by $C_T^\tau F_T$ and the allowed values of the cosine of the weak phase $\cos\phi$ for $Case(a)$ I taken from Table II of [8].

$Case(a)$ II:

Similar to the $Case(a)$ I, using the average width $\frac{\Gamma^{\tau^+}+\Gamma^{\tau^-}}{2}$, the CP asymmetry due to direct CP violation $A_{CP} = \frac{\Gamma^{\tau^+}-\Gamma^{\tau^-}}{\Gamma^{\tau^+}+\Gamma^{\tau^-}}$ and using the combination $K^*(892) + K^*(1410) + K^*(800)$, we get two feasible solutions for $C_T^\tau F_T$ and ϕ as shown in Table 5.2 where here also only the first solution is viable.

Sl.No	$C_T^\tau F_T$	$\cos\phi$	$\lvert\frac{A_T}{A_{SM}}\rvert^2$	$\frac{Int(A_T A_{SM})}{A_{SM}^2}$(cos term)	$\frac{Int(A_T A_{SM})}{A_{SM}^2}$(sin term)
(i)	-0.303	-0.970	0.018	0.097	0.008
(ii)	-1.945	-0.990	0.759	0.638	0.028

Table 5.2: Table showing the NP contribution parameterized by $C_T^\tau F_T$ and the allowed values of the cosine of the weak phase $\cos\phi$ for $Case(a)$ II taken from Table I of [8].

$Case(b)$:

In $Case(a)$ we have assumed that the tensor form factors are constant, but in general the tensor form factor is expected to depend on Q^2 just like the vector and scalar form factors, one way to probe the Q^2 dependence of the tensor form factor is to fix its parameters, such as Q^2 dependent mass and width in the Breit-Wigner form, by comparing with experimental data. Another possible way to probe the Q^2 dependence of the $F_T(Q^2)$ is to use Dirac equations of motion to express the tensor form factors in terms of scalar and vector form factors, we have done a rudimentary analysis in this direction in [19].

5.4 Summary

Babar collaboration has reported an intriguing opposite sign in the integrated decay rate asymmetry $A_{cp}(\tau \rightarrow K_S \pi \nu_\tau)$ than that of SM prediction from the known K^0 - $\bar{K^0}$ mixing. Babar's result deviate from the SM prediction by about 2.7σ. In this work we have shown that NP of scalar type or vector type will not contribute to CP violation in the integrated rate in $\tau \rightarrow K_S \pi \nu_\tau$, and so if the observed CP asymmetry in integrated decay rate in $\tau \rightarrow K_S \pi \nu_\tau$ stand with further experimental test, then only tensor current NP can explain this observation. Using the Breit Wigner forms of the various resonances contributing to the vector and scalar form factors, the $s(K_S\pi)$

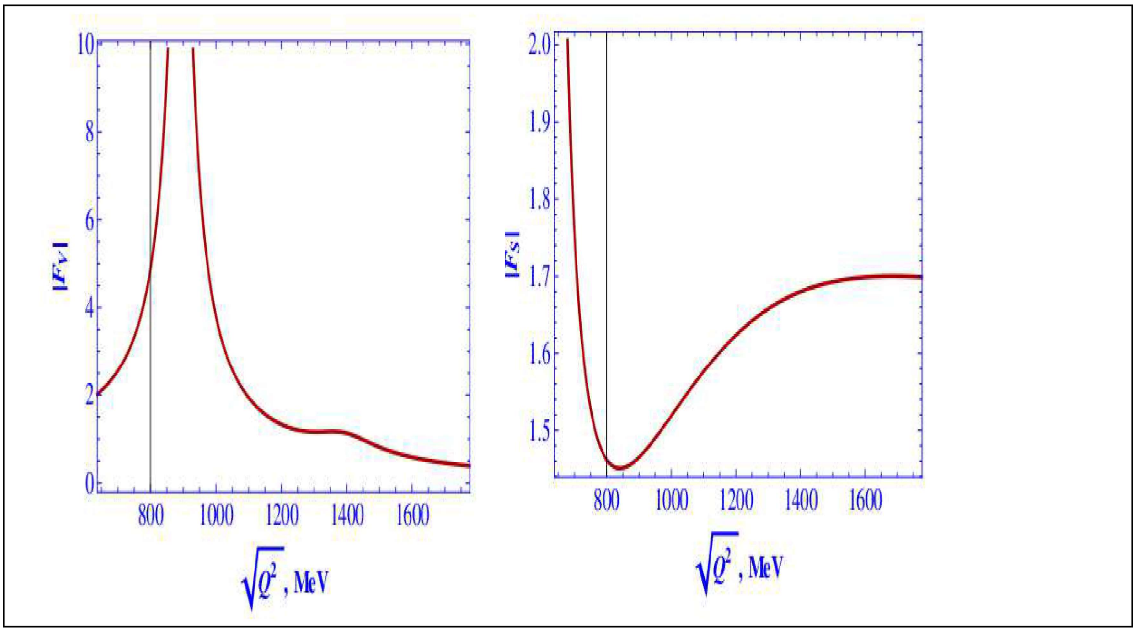

Figure 5.6: Plot of $|F_v|$ as a function of $\sqrt{Q^2}$ (in MeV) with contributions coming from $K^*(892)$ and $K^*(1410)$ (left) and lot of $|F_s|$ as a function of $\sqrt{Q^2}$ (in MeV) with contributions coming from $K^*(800)$ (right) [8].

dependent strong phases were determined and then taking the strong phase of the tensor form factor to be vanishing and its magnitude to be constant, the observables $A_{cp}(\tau \to K_s \pi \nu_\tau)$ and branching fraction were used to determine weak phase and the magnitude of the tensor coupling. In this work we have also given a full angular distributions of the decay rate and CP asymmetry in presence of new tensor current. In future, whether the experimentalist fit the parameters of the $F_T(Q^2)$ from their data, or a more rigurous relation derived using Dirac equations of motion to express $F_T(Q^2)$ in terms of $F_V(Q^2)$ and $F_S(Q^2)$ [19], then we can do a full Q^2 dependence analysis.

Chapter 6

Concluding Remarks

To conclude, flavor measurements in general provide very sensitive ways to probe SM and possible effects from New Physics(NP), and also it can put very strong constraints on the parameters of NP models. Currently there are several measurements showing deviations from SM predictions in the range of (2–4)σ level. In the following paragraphs, we will give a brief chapter wise summary of this thesis.

In chapter I, we have given a brief introduction of the key concepts underpinning the SM of particle physics, like Yang-Mills principle of local guage invariance, spontaneous symmetry breaking and Higgs mechanism.

In chapter II, we have presented a brief review of the flavor physics, especially related to the experimentally observed anomalies by Babar, Belle and LHCb in recent years. Currently there are several measurements showing deviations from SM predictions in the range of (2–4)σ level. In following we will enumerate the most significant ones briefly. $R_K = \frac{Br(\bar{B}\to\bar{K}^{(*)}\mu^+\mu^-)}{Br(\bar{B}\to\bar{K}^{(*)}e^+e^-)}$: The LHCb collaboration announced a 2.6 σ deviation in the measurement of the ratio of the branching fraction of B to $\bar{K}^{(*)}$ and dimuons to that of B to $\bar{K}^{(*)}$ and dielectrons [50]. Muon $(g-2)$: The muon $(g-2)$ measurements has been reported to deviate from the SM expectation by more than 3σ though there still are quite large uncertainties in the SM predictions of this process.

$R(D^{(*)}) = \frac{Br(B\to D^{(*)}\tau\nu_\tau)}{Br(B\to D^{(*)}l\nu_l)}$: BaBar[20], Belle[21] and LHCb[22] measurements of $R(D^{(*)})$ decays with respect to SM predictions for these decays shows about 4σ(if we take the three deviations together). One of the main topic of this thesis is concerned with explaining this anomaly. Also there is the Babar collaboration[4]'s reported anomaly in the measurement of CP asymmetry in $\tau \to K_S\pi\nu_\tau$ in the time integrated decay rate.

Chapter III has focused on the "Possible Hints of Lepton Flavor Universality Violation in R(D$^{(*)}$)". In the following paragraphs we will give a summary of the key results and conclusions[27]. BaBar[20], Belle[21] and LHCb[22] have reported an excess in the measurements of $R(D^*)$, $R(D)$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ than expected from SM, a possible signature of lepton flavor universality violating NP. In this work we have analyzed the implications for these decay modes in a Flipped two Higgs doublet model(2HDM) with enhanced Yukawa coupling of $H^\pm$ to τ lepton. In Table 6.1 we have shown two different values of the parameters $\tan\beta$ and $M_\pm$ that fits at same accuracy:

S.no	$\tan\beta$	$M_\pm$ GeV	$R(D)_{Th}$	$R(D^*)_{Th}$	$Br_{Th}(B\to\tau\nu)$
1	69.97	700	0.348	0.255	1.29×10^{-4}
2	99.95	1000	0.348	0.255	1.29×10^{-4}

Table 6.1: $\chi^2_{min} = 10.95$ and we have restricted the $\tan\beta$ in the range of $100 > \tan\beta > 1$.

From Table 6.1, we see that in the range 1 TeV $\geq M_\pm \geq$ 600 GeV and $100 > \tan\beta >$

1, we have :

$$R(D)_{Th} = 0.348 \pm 0.16 \tag{6.1}$$

$$R(D^*)_{Th} = 0.255 \pm 0.07 \tag{6.2}$$

and

$$Br_{Th}(B \to \tau\nu) = (1.29 \pm 0.89) \times 10^{-4}, \tag{6.3}$$

compared to the combined [Babar,Belle,LHCb][23] experimental values:

$$R(D)_{EXP} = 0.388 \pm 0.047 \tag{6.4}$$

$$R(D^*)_{EXP} = 0.321 \pm 0.021 \tag{6.5}$$

and

$$Br_{EXP}(B \to \tau\nu) = (1.14 \pm 0.27) \times 10^{-4}. \tag{6.6}$$

By adding theoretical and experimental errors in quadrature from Eqs. (6.1,6.2,6.3) and Eqs. (6.4,6.5,6.6), we conclude that our phenomenological model can give results in agreement within 1σ deviation for the combination of $R(D^{(*)})$ and $\mathcal{B}r(B \to \tau\nu_\tau)$ compared to about 4σ deviation from SM for the latest combined[Babar,Belle,LHCb] experimental data for these observables. The same results can be achieved if b quark replaces the τ lepton in a Lepton Specific 2HDM. In that case Yukawa coupling of the leptonic sector will be same as Type-II 2HDM and Yukawa coupling of the quark sector will be same as Lepton Specific 2HDM [29] except the b quark which will have the effective Yukawa coupling same as in Type-II with opposite sign, and the charged Higgs mass may be allowed to be lower than 600 GeV in that case. We also observed that from the form of the Yukawa couplings it is expected that if we require $\eta = -1$ for the

b quark or τ lepton in the 2HDM-II, then also it will fit the three observables at about same acuracy as above model of Flipped/Lepton-Specific 2HDM. This is interesting in a sense that it will be like 2HDM-II but with a wrong sign in either b quark or τ lepton Yukawa coupling, an anomalous SUSY.

In Chapter IV, we have presented a review of the theoretical analysis of CP violation in hadronic τ decays as well the experimental status of the some of the key CP observables in hadronic τ decays.

In Chapter V we have presented another main topic of this thesis "Observed Sign Anomaly in $A_{cp}(\tau \to K_s \pi \nu_\tau)$". In the following paragraphs we will give a summary of the key results and conclusions of our analysis of the contributions coming from tensorial current to this observable [8][19]. Babar collaboration has reported an intriguing opposite sign in the time integrated decay rate asymmetry in $\tau \to K_s \pi \nu_\tau$ than that of SM prediction from the known K^0 - $\bar{K^0}$ mixing. Babar [4]'s result deviate from the SM prediction by about 2.7σ. In [8], we have shown that the CP violation coming from the $K - \bar{K}$ mixing and the direct CP violation in $A_{cp}(\tau \to K_s \pi \nu_\tau)$ can be separated as

$$A_{cp}(\tau \to K_s \pi \nu_\tau) = \frac{A_{cp}^K + A_{cp}^\tau}{1 + A_{cp}^K A_{cp}^\tau} \tag{6.7}$$

where A_{cp}^K is the CP violation coming from the $K - \bar{K}$ mixing and A_{cp}^τ is the direct CP violation coming from NP particle mediated CP violating at lepton and/or quark vertices and τ_τ is the τ life time. Since both A_{cp}^K and A_{cp}^τ are expected to be small, we can safely ignore terms involving the product of the two in the denominator. In [8], using the Breit Wigner forms of the various resonances contributing to the vector and scalar form factors, the $s(K_S\pi)$ (hadronic invariant mass) dependent strong phases were determined. Taking the strong phase of the tensor form factor to be vanishing and its magnitude to be constant, the observables $A_{cp}(\tau \rightarrow K_s\pi\nu_\tau)$ and branching fraction were used to determine weak phase and the magnitude of the tensor coupling. Also we have given a full angular distributions of the decay rate and CP asymmetry in presence of new tensor current of the $\tau \rightarrow K_s\pi\nu_\tau$.

Appendix A

Separating The Mixing And The Direct CP Violation Parts

In presence of New Physics, the experimentally measured CP asymmetry A_{CP}^{Exp} will be result of the combination of CP asymmetry arising from the $K^0 - \bar{K}^0$ mixing and direct CP aymmetry coming from the leptonic current vertex or the quark current vertex. Hence in the presence of both direct and indirect CP violations, we can express the observed decay rate asymmetry as[8]

$$A_{CP}^{Exp} = \frac{\Gamma(\tau^+ \to K_S\pi^+\nu_\tau)\int_{t_1}^{t_2}\Gamma(K^0 \to \pi\pi) - \Gamma(\tau^- \to K_S\pi^-\nu_\tau)\int_{t_1}^{t_2}\Gamma(\bar{K}^0 \to \pi\pi)}{\Gamma(\tau^+ \to K_S\pi^+\nu_\tau)\int_{t_1}^{t_2}\Gamma(K^0 \to \pi\pi) + \Gamma(\tau^- \to K_S\pi^-\nu_\tau)\int_{t_1}^{t_2}\Gamma(\bar{K}^0 \to \pi\pi)} \tag{A.1}$$

Now if we define

$$A_{CP}^{K^0-\bar{K}^0} = \frac{\int_{t_1}^{t_2}[\Gamma(K^0 \to \pi\pi) - \Gamma(\bar{K}^0 \to \pi\pi)]}{\int_{t_1}^{t_2}[\Gamma(K^0 \to \pi\pi) + \Gamma(\bar{K}^0 \to \pi\pi)]} = \frac{\int_{t_1}^{t_2}[\Gamma^{K^0} - \Gamma^{\bar{K}^0}]}{\int_{t_1}^{t_2}[\Gamma K^0 + \Gamma^{\bar{K}^0}]} \tag{A.2}$$

and also $\Gamma^{\tau^+} = \Gamma(\tau^+ \to K_S\pi^+\nu_\tau)$ and $\Gamma^{\tau^-} = \Gamma(\tau^- \to K_S\pi^-\nu_\tau)$ with

$$A_{CP}^{\tau} = \frac{\Gamma^{\tau^+} - \Gamma^{\tau^-}}{\Gamma^{\tau^+} + \Gamma^{\tau^-}} \tag{A.3}$$

then we have

$$A_{CP}^{Exp} = \frac{\Gamma^{\tau^+} \int_{t_1}^{t_2} \Gamma^{K^0} - \Gamma^{\tau^-} \int_{t_1}^{t_2} \Gamma^{\bar{K}^0}}{\Gamma^{\tau^+} \int_{t_1}^{t_2} \Gamma^{K^0} + \Gamma^{\tau^-} \int_{t_1}^{t_2} \Gamma^{\bar{K}^0}}. \tag{A.4}$$

Now we have

$$\Gamma^{\tau^+} \int_{t_1}^{t_2} \Gamma^{K^0} - \Gamma^{\tau^-} \int_{t_1}^{t_2} \Gamma^{\bar{K}^0} = 2\{\tfrac{\Gamma^{\tau^+}+\Gamma^{\tau^-}}{2} \tfrac{\int_{t_1}^{t_2}[\Gamma^{K^0}-\Gamma^{\bar{K}^0}]}{2} + \tfrac{\Gamma^{\tau^+}-\Gamma^{\tau^-}}{2} \tfrac{\int_{t_1}^{t_2}[\Gamma^{K^0}+\Gamma^{\bar{K}^0}]}{2}\}$$

$$= \frac{1}{2}(\Gamma^{\tau^+} + \Gamma^{\tau^-}) \int_{t_1}^{t_2} [\Gamma^{K^0} + \Gamma^{\bar{K}^0}][A_{CP}^{K^0} + A_{CP}^{\tau}] \tag{A.5}$$

and similarly the sum is given as

$$\Gamma^{\tau^+} \int_{t_1}^{t_2} \Gamma^{K^0} + \Gamma^{\tau^-} \int_{t_1}^{t_2} \Gamma^{\bar{K}^0} = \frac{1}{2}(\Gamma^{\tau^+} + \Gamma^{\tau^-}) \int_{t_1}^{t_2} [\Gamma^{K^0} + \Gamma^{\bar{K}^0}][1 + A_{CP}^{K^0} A_{CP}^{\tau}]. \tag{A.6}$$

Therefore, the observed asymmetry for the decay rate can be expressed as

$$A_{CP}^{Exp} = \frac{A_{CP}^{K^0} + A_{CP}^{\tau}}{1 + A_{CP}^{K^0} A_{CP}^{\tau}}. \tag{A.7}$$

In the cases where the NP contribution to A_{CP}^{τ} are comparable or smaller than SM part $A_{CP}^{K^0}$, we can neglect the product term in the denominator and take the above equation as

$$A_{CP}^{Exp} \approx A_{CP}^{K^0} + A_{CP}^{\tau}. \tag{A.8}$$

Then the direct CP violation and CP violation due to oscillation part get decoupled.

Appendix B

Definitions of Helicity Amplitudes, Coordinate Frames, Kinematics And Polarization Vectors

Coordinate Frames, Kinematics And Polarization Vectors.

We choose the coordinates such that, in the rest frame of decaying B meson, the D* momentum is along the z-axis and the charged lepton momentum is in the x-z plane with positive x-component. Then the four momentum is given by

$$p_B^\mu = (m_B, 0, 0, 0),\ p_{D^*}^\mu = (E_{D^*}, 0, 0, |\vec{p}_{D^*}|),\ q^\mu = (q^0, 0, 0, -|\vec{p}_{D^*}|) \tag{B.1}$$

where

$$E_{D*} = \frac{m_B^2 + m_{D^*}^2 - q^2}{2m_B}, \quad q^0 = \frac{m_B^2 - m_{D^*}^2 + q^2}{2m_B}, \quad |\vec{p}_{D^*}| = \frac{\sqrt{Q_+ Q_-}}{2m_B} \tag{B.2}$$

with

$$Q_{\pm} = (m_B \pm m_{D^*})^2 - q^2. \tag{B.3}$$

The polarization vectors for the virtual W in the rest frame of the B meson are given by

$$\epsilon(W,\pm)^\mu = \mp\frac{1}{\sqrt{2}}(0,1,\mp i,0),\ \epsilon(W,0)^\mu = \frac{1}{\sqrt{q^2}}(|\vec{p}_{D^*}|,0,0,-q^0),\ \epsilon(W,s)^\mu = \frac{1}{\sqrt{q^2}}q^\mu. \tag{B.4}$$

and polarizations of the D^* meson in the rest frame of the B meson are given by

$$\epsilon(D^*,\pm)^\mu = \mp\frac{1}{\sqrt{2}}(0,1,\mp i,0), \quad \epsilon(D^*,0)^\mu = \frac{1}{m_{D^*}}(|\vec{p}_{D^*}|,0,0,E_{D^*}). \tag{B.5}$$

In this frame the D^* meson is moving along the positive z-axis, whereas the virtual W is moving in the negative z-axis.
In the rest frame of the virtual W, we have

$$p_B^\mu = (E_B,0,0,|\vec{p}_B|), \qquad q^\mu = (\sqrt{q^2},0,0,0) \tag{B.6}$$

and

$$p_l^\mu = (E_l, |\vec{p}_l|\sin\theta_l, 0, |\vec{p}_l|\cos\theta_l),\ p_\nu^\mu = (|\vec{p}_l|, -|\vec{p}_l|\sin\theta_l, 0, -|\vec{p}_l|\cos\theta_l) \tag{B.7}$$

where

$$E_B = \frac{m_B^2 - m_{D^*}^2 + q^2}{2\sqrt{q^2}},\ p_B = \frac{1}{2\sqrt{q^2}}\sqrt{Q_+Q_-},\ E_l = \frac{q^2+m_l^2}{2\sqrt{q^2}},\ |\vec{p}_l| = \frac{q^2-m_l^2}{2\sqrt{q^2}} \tag{B.8}$$

where θ_l is the angle measured between the lepton l and the D^* meson in the rest frame of the virtual W and can be expressed as

$$\cos\theta_l = \frac{(q^2+m_l^2)(m_B^2 - m_{D^*}^2 + q^2) - 4m_B^2 q^2 x}{(q^2-m_l^2)\sqrt{Q_+Q_-}} \tag{B.9}$$

with

$$x = \frac{p_l \cdot p_B}{m_B^2} \tag{B.10}$$

and in the rest frame of the virtual W, the polarization of the virtual W is given as

$$\epsilon(W^*, \pm)^\mu = \mp\frac{1}{\sqrt{2}}(0, 1, \mp i, 0),\ \epsilon(W^*, 0)^\mu = (0, 0, 0, -1),\ \epsilon(W^*, s)^\mu = (1, 0, 0, 0). \tag{B.11}$$

Helicity Amplitudes And Evaluation of The Leptonic Matrix Elements.

In general, the decay $B \to D^{(*)} l \nu_l$ can be expressed as

$$M(B \to D^{(*)} l \nu_l) = \frac{G_F V_{cb}}{\sqrt{2}} L_\mu H^\mu \tag{B.12}$$

by inserting the completeness relation [67][1]

$$\sum_{m,n} \epsilon_\mu(m) \epsilon_\nu^*(n) g_{mn} = g_{\mu\nu} \tag{B.13}$$

where $g_{mn} = diag(1, -1, -1, -1)$ and so we can express equation B.12 as

$$M(B \to D^{(*)} l \nu_l) = \frac{G_F V_{cb}}{\sqrt{2}} \sum_{m,n} L(m) H(n) g_{mn} \tag{B.14}$$

where

$$H(n)_{\lambda_{D^*}, \lambda_W} = \epsilon_\mu^*(n, \lambda_W) \langle D^*(\lambda_{D^*}) | J_{cb}^\mu | B \rangle \tag{B.15}$$

[1]intuitively this can be guessed for the Virtual $W^*(M_{W^*}^2 = q^2)$ from the relation $\sum_{m=0}^{2} \epsilon_\mu(m) \epsilon_\nu^*(m) = -g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}$ as $\frac{q_\mu}{\sqrt{q^2}}$ act as the $\epsilon_\mu(s)$, the scalar component of the virtual W^* polarization in the virtual W^* rest frame.

and

$$L(m)_{\lambda_l,\lambda_W} = \epsilon^*_\mu(m,\lambda_W)\langle l(\lambda_l)\nu_l(\lambda_l)|J^\mu_{l\nu_l}|0\rangle \tag{B.16}$$

with λ_l and λ_W as the helicities of lepton l and virtual W respectively.

Now the hadronic matrix elements can be evaluated in terms of the form factors, as done in the sections 3.7 and 3.7, so here we will give the method of evaluating the leptonic matrix elements following the derivation given in [67]. In the chiral representation, the Dirac matrices are given by

$$\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu_+ \\ \sigma^\mu_- & 0 \end{pmatrix}, \qquad \gamma^5 = \begin{pmatrix} -1 & 0 \\ 0 & 1 \end{pmatrix} \tag{B.17}$$

with $\sigma^\mu = (1, \pm\sigma_i)$ where σ_i are the Pauli matrices. The four component spinor can be expressed in terms of the two component spinor as

$$u(p,\lambda) = \begin{pmatrix} u(p,\lambda)_- \\ u(p,\lambda)_+ \end{pmatrix}, \qquad v(p,\lambda) = \begin{pmatrix} v(p,\lambda)_- \\ v(p,\lambda)_+ \end{pmatrix} \tag{B.18}$$

where the λ denotes the helicity of the particle. Now the two component spinors can be expressed in terms of the helicity eigenspinors as

$$u(p,\lambda)_\pm = \omega_\pm\chi(\theta)_+, \qquad v(p,\lambda)_\pm = \pm\omega_\mp\chi(\theta)_- \tag{B.19}$$

where the chirality conserving/flipping factor is given by $\omega_\pm = \sqrt{E \pm p}$ respectively, and the helicity eigenspinors are given by

$$\chi(\theta)_+ = \begin{pmatrix} \cos\frac{1}{2}\theta \\ \sin\frac{1}{2}\theta e^{i\phi} \end{pmatrix}, \qquad \chi(\theta)_- = \begin{pmatrix} -\sin\frac{1}{2}\theta e^{-i\phi} \\ \cos\frac{1}{2}\theta \end{pmatrix} \tag{B.20}$$

which dsecribe either particles of helicity $\pm\frac{1}{2}$ respectively or anti-particles of helicity $\mp\frac{1}{2}$ respectively having an arbitary four-momentum given by

$$p^\mu = (E, p\sin\theta\cos\phi, p\sin\theta\sin\phi, p\cos\theta) \tag{B.21}$$

where $0 \leq \theta \leq \pi$ and $0 \leq \phi \leq 2\pi$, but in our choice of coordinate system, $\phi = 0$, so the matrix elements are independent of the angle ϕ.

So then in the rest frame of the virtual W, the leptonic matrix elements $L(m)_{\lambda_l,\lambda_W}$ given in Eqs(C.16) can be evaluated as

$$\begin{aligned} L(m)_{\lambda_l,\lambda_W} &= \epsilon^*_\mu(m,\lambda_W)\bar{u}(p_l,\lambda_l)\gamma^\mu(1-\gamma^5)v(p_\nu,+) \\ &= -2\omega(p_l)_-\omega(p_\nu)_+\epsilon^*_\mu(m,\lambda_W)\chi(\theta)^\dagger_{\lambda_l}\sigma^\mu_-\chi(\theta)_{\lambda_\nu=-} \end{aligned} \tag{B.22}$$

then using Eqs(C.20) with $\phi = 0$ we have

$$\chi(\theta)^\dagger_-\sigma^\mu_-\chi(\theta)_- = (0, -\cos\theta, -i, \sin\theta) \tag{B.23}$$

and

$$\chi(\theta)^\dagger_+\sigma^\mu_-\chi(\theta)_- = (0, -\sin\theta, 0, -\cos\theta). \tag{B.24}$$

Then using the form of the virtual W polarizations from the Eqs.(C.11) we have

$$L(m)_{-,\pm} = 2\sqrt{q^2}vd_\pm,\ L(m)_{-,0} = -2\sqrt{q^2}vd_0,\ L(m)_{-,s} = 0 \tag{B.25}$$

and

$$L(m)_{+,\pm} = \pm 2m_l v d_0,\ L(m)_{+,0} = -2m_l v(d_+ - d_-),\ L(m)_{+,s} = -2m_l v \tag{B.26}$$

with

$$v = \sqrt{1 - \frac{m_l^2}{q^2}},\ d_\pm = \frac{1 \pm \cos\theta}{\sqrt{2}},\ d_0 = \sin\theta \tag{B.27}$$

where θ here is the angle made by the charged lepton l with the z-axis (direction of the decaying B meson) in the rest frame of the virtual W.

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