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SELECTED LOW ENERGY PROBES FOR NEW PHYSICS

By

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"The effort to understand the universe is one of the very few things that lifts human life a little above the level of farce, and gives it some of the grace of tragedy".

Steven Weinberg

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Abstract

In this work we select two interesting methods in the intensity frontier of particle physics, as low energy probes of New Physics. First, we address a crucial unknown property of the neutrino, its Majorana or Dirac nature. The Majorana nature of the neutrino can be established via the observation of lepton number violating processes. We propose a new strategy for detecting the Charge and Parity (CP) violating phases and the effective mass of muon Majorana neutrinos by measuring observables associated with neutrino-antineutrino oscillations in $\pi^{\pm}$ decays. Within the generic framework of quantum field theory, we compute the non-factorizable probability for producing a pair of same-charged muons in $\pi^{\pm}$ decays as a distinctive signature of $\nu_{\mu}-\bar{\nu_{\mu}}$ oscillations. We show that an intense neutrino beam through a long baseline experiment is favored for probing the Majorana phases. Using the neutrino-antineutrino oscillation probability reported by MINOS collaboration, a new stringent bound on the effective muon-neutrino mass is derived.

Secondly, we aim to constrain New Physics with allowed processes within the Standard Model (SM) but measured with very good precision, enough to test for its accuracy. We choose an interesting set of observables, not yet exploited, and within the range of charm physics, the D meson leptonic and semileptonic decays. Latest Lattice results on D form factors evaluation from first principles show that the standard model (SM) branching ratios prediction for the leptonic $D_s \to \ell\nu_\ell$ decays and the semileptonic SM branching ratios of the D^0 and D^+ meson decays are in good agreement with the world average experimental measurements. It is possible to disprove New Physics hypothesis or find bounds over several models beyond the SM. Using the observed leptonic and semileptonic branching ratios for the D meson decays, we performed a combined analysis to constrain non standard interactions

which mediate the $c\bar{s} \to l\bar{\nu}$ transition. This is done either by a model independent way through the corresponding Wilson coefficients or in a model dependent way by finding the respective bounds over the relevant parameters for some models beyond the standard model. In particular, we obtain bounds for the Two Higgs Doublet Model Type-II and Type III, the Left-Right model, the Minimal Supersymmetric Standard Model with explicit R-Parity ($R = (-1)^{3B+L+2S}$, where B, L and S are the baryon number, lepton number and particle spin respectively) violation and Leptoquarks. Finally, we estimate the transverse polarization of the lepton in the D^0 decay and we found it can be as high as $P_T = 0.23$.

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I specially dedicate this work to my daughter Lilieth.

Introduction

The physics that beautifully describes the interactions and fundamental properties of the particles at subatomic scales, the Standard Model, is extraordinary precise and accurate, with no significant deviations from it. Last year in fact we witnessed the discovery of one of the last pieces, the Higgs scalar boson [1, 2]. The Standard Model (SM) is a consistent, finite, and within our technical limitations, measurable theory. The SM describes the strong, weak and electromagnetic interactions, specified by the gauge theory $SU(3)_c \times SU(2)_L \times U(1)_Y$. The energy range in which it has been tested is wide: from atomic processes to the Z-pole and beyond. At least no direct observation of fundamental new particles has been done up to energy scales of the order of TeV, almost ten times larger than the Z boson mass.

Despite its success, there are strong observations and crucial theoretical reasons to expect a more fundamental theory. It excludes in its description the gravitational interaction, which is known to become relevant at extremely high energies, which are for now unreachable in artificial particle accelerators. The cosmological observations show us compeling evidence of a large amount of dark matter, much larger than the common barionic matter, and a visible asymmetry between matter and antimatter that can not be explained within the SM. It seems that many of the phenomena occuring at large scales that would have affected the very begining of our universe, where subatomic particles interactions become relevant are beyond the SM description. Even the observed changing flavor of neutrinos (neutrino oscillations) requieres a minimum extension of the SM to include massive neutrinos. Many other issues may be addressed that have nule or insatisfactory explanations within the SM, as the quadratically divergent quantum corrections to the Higgs boson mass, the origin of parity-violation in weak interactions; the quantization of electric charge; and values of the input parameters, such as the fermion masses whose spectrum spans eleven orders of magnitude [3].

In order to search for New Physics, physics beyond the Standard Model (BSM), physicists explore three frontiers [3]: the high energy frontier, now explored by the CERN Large Hadron Collider; the Cosmological frontier, including probes of the cosmic microwave background and large scale structure as well as indirect astrophysical detection of dark matter; and the Intensity or Precision frontier, from atomic processes to high energy physics. We are interested in the latest, the intensity or precision frontier. There are a huge amount of measured observables with extraordinary precision in literature one can use to either test BSM physics or set the corresponding limits in which the SM is without refutation accurate. We aim to provide some methods to follow this direction. The intensity frontier is a low energy probe and indirect search of New Physics. One could search for forbidden processes in the SM, or highly supressed, such as electric dipole moments, lepton number violation, flavor changing neutral processes, etc. Another possibility is to analyse measured quantities allowed within the SM but with sufficient efficiency to test non standard interactions (NSI). Furthermore, neutrino experiments, which have been continuously reaching an extraordinary precision and may be used too as a probe for New Physics.

In this work we select two interesting methods in the intensity frontier to search for New Physics. First, we address a crucial unknown property of the neutrino, its nature. No elementary fermion has been known to be its own antiparticle, *i.e. a Majorana fermion.* The neutrino can either be a Majorana fermion or a Dirac fermion. Low energy experiments, such as neutrinoless double beta decays search for this answer. If neutrinoless double beta decay is observed it necessarily implies lepton number violation, which in the SM does not occur. The Majorana nature of the neutrino can be established via the observation of lepton number violating processes. The parameter characterizing the rate of such transitions involves a combination of neutrino masses, the mixing angles and phases. A direct observation of its effective Majorana mass would be clearly useful. Neutrinoless double beta decay experiments will help set the effective mass for the electron neutrino but can not access the CP assymmetry. Neutrino antineutrino oscillations, although merely unreachable at the time, and starting to be addressed, is a lepton number violating process that is sensitive to the Majorana phases, and the CP asymmetry. Here, we propose a mechanism, based in neutrino antineutrino oscillations, which would allow to derive a strong bound on the effective Majorana mass of the muon neutrino [4]. This is fully described and discussed in chapter one.

Secondly, we aim to constrain New Physics with allowed processes within the SM but measured with very good precision, enough to test for its accuracy. We choose an interesting set of observables, not yet exploited, and within the range of charm physics, the D meson leptonic and semileptonic decays [5]. The hadron physics is much more subtle than lepton interactions. Hadrons are composite particles of fundamental fermions, quarks. The quarks have an extra quantum number as compared with the leptons, the color charge, which means they interact strongly, described by the QCD (quantum chromodynamics) gauge symmetry. At low energies, in the range where the D mesons are produced, the strong interaction among the quarks is not perturbative, in fact they show confinement. One may parameterize our ignorance of the detailed structure of the strong interactions with *form factors*. The SM predictions for the D meson leptonic and semileptonic decays relies on the lattice QCD estimates of the form factors, which has reached an extraordinary precision, and appear to be in agreement with the world average experimental measurements [6], allowing us to disprove New Physics hypothesis or find restrictive bounds over several models beyond the SM. In chapter two, we fully describe a method to constrain NSIs combining the leptonic and semileptonic decays of the D meson.

Finally we give our general perspectives and conclusions in the last chapter of this thesis.

Chapter 1

$\Delta L = 2$ neutrino oscillations

1.1 Introduction

Neutrinos are massive particles, but at present we ignore the hierarchy of their mass spectrum. Measuring the absolute values of these masses would provide important clues to establish the mechanism that gives rise to neutrino masses. Currently, only indirect constraints from experiments can be derived on neutrino masses by means of reasonable assumptions regarding the neutrino mixings and phases. In this way, neutrino mass differences measured in oscillation experiments combined with searches in tritium beta decays and in neutrinoless nuclear double beta decays indicate that all neutrino masses are below the eV scale [7]. On the other hand, detecting CP violating phases in the lepton sector remains one of the most challenging problems in the study of neutrino mixing. In the basis of diagonal charged lepton mass matrix, the neutrino mass matrix m_ν can be written in flavor basis as $m_\nu = U^* m_\nu^{diag} U^\dagger$, where U is the Pontecorvo-Maki-Nagakawa-Sakata (PMNS) mixing matrix, which can be written as $U = V \cdot diag\,(1, e^{i\alpha_2}, e^{i\alpha_3})$. Here α_i are the Majorana phases and the mixing matrix V can be parameterized by one Dirac phase and three angles: solar angle θ_{12}, atmospheric angle θ_{23}, and Chooz angle θ_{13}.

As is well known, the Majorana nature of neutrinos can be established via the observation of $\Delta L = 2$ processes [8, 9],[10]. The parameter characterizing the rate of such transitions, the effective neutrino mass $\langle m_{ll}\rangle \equiv \sum_i U_{li}^2 m_{\nu_i}$, involves a combination of neutrino masses, mixings and phases. Although it is clear that a combination of neutrino oscillation and nuclear beta decay experiments currently provide strong constraints on neutrino masses, having direct constraints on effective Majorana masses would be useful to have a more complete picture of Majorana neutrinos. Actually, it turns out that the only way to access the values of Majorana phases is through observables associated to $\Delta L = 2$ transitions [7].

Nevertheless, measurements of the effective electron-neutrino mass in the neutrinoless double beta decay ($0\nu\beta\beta$) experiments can not restrict the two Majorana CP violating phases present in the PMNS mixing matrix. The effective electron-neutrino mass $\langle m_{ee}\rangle$ is given by $|\langle m_{ee}\rangle| = \left|\sum_i U_{ei}^2 m_{\nu_i}\right|$. This effective mass parameter depends on the angles θ_{12} and θ_{13}, the neutrino masses m_{ν_i}, Dirac CP phase, and Majorana phases α_i. There are several studies on using the results of ($0\nu\beta\beta$) together with the new data from terrestrial and astrophysical observation in order to restrict the Majorana neutrino CP violating phases[7, 11, 12]. However, this analysis is model dependent and quite sensitive to the ansatz of the neutrino mass spectrum: quasi-degenerate, normal or inverted hierarchies. In this respect, it is not possible to measure the Majorana neutrino CP phases from ($0\nu\beta\beta$) experiment. This may be expected since in ($0\nu\beta\beta$) one measures the lifetime of the decay of two neutrons in a nucleus into two protons and two electrons, which is a CP conserving quantity [13].

On the other hand, direct bounds on other effective neutrino mass parameters $\langle m_{ll}\rangle$ from present experimental data are very poor. Currently, the strongest bound for the muon-neutrino case from the $K^+ \to \pi^-\mu^+\mu^+$ branching fraction [14] is only $|\langle m_{\mu\mu}\rangle| \leq 0.04$ TeV [15], which leads to a negligible constraint on the neutrino masses and CP violating phases.

Therefore, it is commonly believed that direct bounds from other $\Delta L = 2$ decays are only of academic interest [16]; although it is possible to get strong bounds on effective Majorana masses by combining oscillation, kinematical and cosmological data with additional assumptions about Majorana phases, the Majorana nature of neutrinos can be manifest only by observing $\Delta L = 2$ processes. Some attempts to detect CP violation based on the difference between oscillation probabilities of neutrinos and antineutrinos can be found in Ref.[17]. Other proposals aiming to gain access to CP-violating phases of Majorana neutrinos using neutrino-antineutrino oscillations were first discussed in [18, 19, 20, 21, 22, 23, 24, 25].

Here we propose a mechanism [4], based on neutrino-antineutrino oscillation, which would allow to derive a strong bound on the effective Majorana mass of the muon-neutrino $\langle m_{\mu\mu} \rangle$. In addition, it provides a method for detecting the Majorana neutrino CP violating phases through measuring the CP asymmetry of the $\pi^{\pm}$ decay where neutrino-antineutrino oscillation takes place. Using the preliminary bound on the neutrino-antineutrino oscillation probability reported by the MINOS Collaboration [26], we derive a bound on $\langle m_{\mu\mu} \rangle$ which improves existing bounds by several orders of magnitude.

It is worth noting that the probability of a process associated to neutrino oscillation is usually assumed to be factorized into three independent parts: the production process, the oscillation probability and the detection cross section. Since neutrinos carry angular momentum, this may introduce a correlation in the amplitude between their production and detection vertices which eventually can invalidate the factorization hypothesis. In Ref. [27], a generic framework based on quantum field theory was proposed to derive an expression for the time-dependent S-matrix amplitude of neutrino oscillations which avoids the factorization approximation. In order to further test this factorization hypothesis, we use the S-matrix method to compute the time evolution amplitude that describes neutrino-antineutrino conversion from the production to the detection vertices [4]. In section 1.2.1

we briefly review some aspects of neutrino mixing relevant to our work. Secondly we derive expressions for the associated time-dependent CP asymmetries (see section 1.2.2) in the neutrino antineutrino conversion. In section 1.2.3 we use current bounds on the probability for neutrino-antineutrino conversion obtained by the MINOS colaboration [26] in order to provide a bound on the effective Majorana mass of the muon neutrino.

In section 1.3 as an illustrative exercise, we describe the mechanism that would allow us to constrain lepton number violating interactions. In this case we interpret the observation of the final states, same-charged muons at the production and detection of neutrinos, as a result of lepton number violating interactions in pion decays at the neutrino source. Such interactions appear for example in SUSY models with R-Parity violating terms and leptoquark models. In particular in the MSSM without R-Parity conservation [28, 29], the radiative contributions are proportional to the R parity couplings λ and λ', which in general are complex. We attempt to impose constraints over these couplings λ and λ' from the current bound on neutrino-antineutrino transitions obtained by the MINOS collaboration [26].

1.2 Majorana massive states in neutrino-antineutrino oscillations

1.2.1 Neutrino mixing

Let us review briefly some properties of neutrino particles. From a phenomenological point of view, the problem of neutrino mass is strictly related to lepton number conservation: both individual Le,Lμ,Lτ) and total ($\Delta L = 0$), and the Dirac or Majorana nature of neutrinos themselves [24]. In fact, if neutrinos were massless, the conservation of all lepton numbers is allowed, according to the Glashow-Weinberg-Salam theory [30]-[33], due to the invariance of the electroweak lagrangian for un arbitrary (global) phase transformation of the matter field.

Instead, if neutrinos have a non-vanishing mass and are Dirac particles, we can introduce a mass term if $\nu_R^{(i)}$ are added as singlets of the Standard Model $SU(3)_c \times SU(2)_L \times U(1)_Y$, such that $\nu_R^{(i)} \neq \nu_L^{(j)c}$ if $i \neq j$, this is,

$$L_{mass} = -\bar{\nu}_{\alpha_R} m_{\alpha\beta} \nu_{\beta_L} + h.c. \tag{1.1}$$

which is induced by the Higgs scalar when the electroweak symmetry is spontaneously broken, through the Yukawa coupling $L_Y = f_\nu \bar{\nu}_R \phi l_L$. The subscript L denotes chiral left-handed projection and R denotes chiral right-handed projection, i.e. the projectors given by $P_L = (1-\gamma_5)/2$, $P_R = (1+\gamma_5)/2$. Here, $\alpha, \beta = e, \mu, \tau$ and $m_{\alpha\beta}$ is in general, a complex non diagonal 3×3 matrix. Therefore individual lepton number is no longer conserved, while the total lepton number is still conserved. To diagonalize $m_{\alpha\beta}$, we apply a bi-unitary transformation to the fields, $\nu_{\alpha_L} = \sum_{i=1}^{3} U_{\alpha i}^{*} |\nu_{i_L}>$ and $\nu_{\alpha_R} = \sum_{i=3}^{3} V_{\alpha i} |\nu_{i_R}>$, such that

$$L_{mass} = -\bar{\nu}_R^i m_{i_{diag}} \nu_L^i + \text{h.c.} \tag{1.2}$$

where $V^\dagger m U^* = m_{diag}$ and h.c. is the abbreviation for hermitian conjugate. Note that if we include right singlets, this number is completely arbitrary. As a consequence the number of free parameters in the SM will depend on the number of massive neutrinos and its nature.

However, if neutrinos are massive Majorana particles, nor individual or total lepton number is conserved. The particle states satisfy the Majorana condition: $\psi^c = \pm\psi$, which means that each particle is its own antiparticle. If there are no right handed neutrinos in nature we can only construct a Majorana mass term, $\bar{\nu_L^c} \nu_L$, and since $\bar{\nu_L^c} \nu_L$ is a triplet under the $SU(2)$ symmetry, the simplest mass term is $L = \frac{G}{2M} \bar{\nu_L^c} \nu_L \phi\phi$ and the neutrino mass is given by $m_\mu = G < \phi_0^2 > /M$. Note that this Lagrangian is no longer renormalizable and M is an effective mass. For three generations the mass term will be given by $L_{mass} = -\frac{1}{2}\nu_L^{\alpha T} C^{-1} m_{\alpha\beta} \nu_L^\beta + h.c.$.

We diagonalize the former matrix by transforming the neutrino fields as $\nu_L^\alpha = U^*_{\alpha i}\nu_L^i$ so that

$$L_{mass} = -\frac{1}{2}\nu_L^{iT} C^{-1} U_{i\alpha} m_{\alpha\beta} U^*_{\beta j} \nu_L^j + h.c. = -\frac{1}{2}\nu_L^{iT} C^{-1} m_i \nu_L^j + h.c. \tag{1.3}$$

i.e., $U^T m U^* = m_{diag}$.

The phenomenological consequences of eq. 1.2 or 1.3 are precisely neutrino flavour oscillations. But, furthermore, many other processes involving charged leptons, which violate individual lepton number conservation (for Dirac neutrinos) or total lepton number (for Majorana neutrinos), would be allowed.

The neutrino coupling to charged leptons and the W-boson is given within the Standard Model by the following interaction Lagrangian,

$$L_{int} = -\frac{g}{\sqrt{2}} W_\mu^+ \sum_{\alpha i} \nu_{iL} \gamma^\mu U^*_{\alpha i} \ell_{\alpha L} - \frac{g}{\sqrt{2}} W_\mu^- \sum_{\alpha i} \ell_{\alpha L} \gamma^\mu U_{\alpha i} \nu_{iL} \tag{1.4}$$

for massive Dirac or Majorana neutrinos. Here, g is the semiweak coupling constant and the parameter α runs over the charged lepton flavors e, μ, τ. The amplitude of the weak interaction process will be proportional to the mixing matrix $U^*_{\alpha i}$, i.e. $Amp(W^+ \longrightarrow l_\alpha^+ + \nu_i) \propto U^*_{\alpha i}$.

For the case of 3 generations or families, the matrix U of order 3×3 has 3 mixing angles (real parameters) and one complex phase if neutrinos are Dirac particles (similarly for the quark sector), or 3 if neutrinos are Majorana. These phases are CP (Charge-Parity) violating. CP violation however has not been observed in the leptonic sector, and neutrino experiments are expected to measure its size. This matrix is represented by $U = Ue^{i\beta'}e^{i\gamma'}$. For Dirac neutrinos, β', γ' can be absorbed by a phase rotation. If neutrinos are Majorana, α', β' can

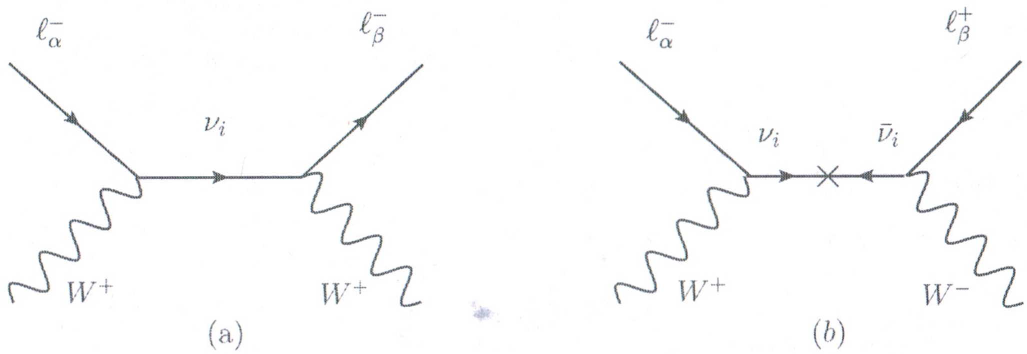

Figure 1.1: Feynman diagrams for (a) neutrino-neutrino and (b) neutrino-antineutrino oscillations. Note that in the second diagram arises a chirality flip in the neutrino propagator which is proportional to the mass m_i of the Majorana neutrino $\nu = \nu^c$. The initial and final neutrino flavor eigenstates are produced and detected via the weak charged-current interactions.

not be absorbed since the mass term takes the form $\nu^{iT} U^T m U \nu^j$. The parametrization of the matrix U is the same as in the quark sector for the Cabbibo-Kobayashi-Maskawa (CKM) matrix, and is given by [7]

$$U_{PMNS} = \begin{array}{c} \\ \nu_e \\ \nu_\mu \\ \nu_\tau \end{array} \begin{array}{c} \begin{array}{ccc} \nu_1 & \nu_2 & \nu_3 \end{array} \\ \left(\begin{array}{ccc} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\varphi} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\varphi} & -c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\varphi} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\varphi} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\varphi} & c_{23}c_{13} \end{array} \right) \end{array}$$
$$\times diag(e^{\frac{\alpha_1}{2}}, e^{\frac{\alpha_2}{2}}, 1) \tag{1.5}$$

where $c_{ij} = \cos\theta_{ij}$ y $s_{ij} = \sin\theta_{ij}$ are the mixing angles, φ is the CP violating phase, α_1, α_2 are the Majorana phases. This matrix is known as the the Pontecorvo-Maki-Nakagawa-Sakata matrix. The mixing in the leptonic sector is much bigger than in the quark sector, i.e. $s_{13} << s_{12} < s_{23} \sim 1$.

The neutrino oscillation may be viewed as the process $l_\alpha^- W^+ \to \nu \to l_\beta^- W^+$ in which the intermediate state neutrino propagates a macroscopic distance L. This process is shown schematically in fig. 1.1. The intermediate-state neutrino is a coherent superposition of the mass eigenstates, so it can be in any of the mass eigenstates ν_i . Thus, the amplitude for this lepton-number conserving process (A_L) may be written as [13]

$$A_L =< W^+ l_\alpha^- | H_{int} | \nu >< \nu | H_{int} | W^+ l_{beta}^- > \tag{1.6}$$

where H_{int} is the interaction Hamiltonian. A_L is proportional to $U_{\alpha i}^* Prop(\nu_i) U_{\beta i}$ and the neutrino flavor oscillation is described essentially by the product of the *factorized amplitudes* $Amp(\nu_\alpha \longrightarrow \nu_\beta) = \sum_i U_{\alpha i}^* Prop(\nu_i) U_{\beta i}$. Here, $Prop(\nu_i)$ is the neutrino propagation through the medium (vacuum or matter). To find the neutrino propagation in vacuum for example, one solves for the usual evolution operator U^{itH} assuming that the neutrinos propagate coherently with definite momentum as plane waves, which at first approximation describes with good precision the neutrino oscillation measured observables. With the former assumptions the probability to find a specific ν_β flavor, $P(\nu_\alpha \longrightarrow \nu_\beta)$, is given by

$$\begin{aligned} P(\nu_\alpha^\pm \longrightarrow \nu_\beta^\pm) \simeq \delta_{\alpha\beta} - 4 \sum_{i>j} \Re(U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^*) \sin(1.27 \triangle m_{ij}^2 \frac{L}{E}) \\ \pm 2 \sum_{i>j} \Im(U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^*) \sin(2.54 \triangle m_{ij}^2 \frac{L}{E}) \end{aligned} \tag{1.7}$$

for relativistic neutrinos. The constant $\triangle m_{ij}^2 = m_i^2 - m_j^2$ is the difference of the square masses of the neutrinos in eV^2, L is given in kilometers and E in GeV. The state ν_α^+ is a neutrino and ν_α^- is an antineutrino.

The best fit values from accelerator, reactor, atmospheric, and solar neutrino experiments

[7] at $\pm 1\sigma$ for the 3-neutrino oscillation parameters are:

$$\begin{aligned}\Delta m_{12}^2 &= 7.58^{+0.22}_{-0.26} \times 10^{-5} eV^2 \\ |\Delta m_{31}^2| &= -2.37 \pm 0.15^{+0.43}_{-0.46} \times 10^{-3} eV^2 \\ \sin^2\theta_{12} &= 0.306^{+0.12}_{-0.09} \\ \sin^2\theta_{23} &= 0.42^{+0.08}_{-0.03} \\ \sin^2\theta_{13} &= 0.021^{+0.007}_{-0.008}\end{aligned} \tag{1.8}$$

The hierarchy of the masses, normal $m_1 << m_2 < m_3$ or inverted $m_3 << m_1 < m_2$ is still unknown, as well as the absolute values of the neutrino masses. The cosmological data however is used to obtain an upper limit in the total sum of the neutrino masses and gives $\sum_j m_j \lesssim (0.3 - 1.3)$eV at 90% confidence level.

On the other hand, a qualitatively different process is neutrino-antineutrino oscillation. This can be described by $l_\alpha^+ W^- \to \nu \to l_\beta^- W^+$, where the intermediate neutrino travels a macroscopic distance L. This process is shown in fig. 1.1. Unlike ordinary flavor oscillations, this process can only occur if lepton number is no longer a good quantum number. This is exactly the case if the neutrinos have non-vanishing Majorana masses, which also implies that the neutrino mass eigenstates are Majorana particles. The amplitude will be given by

$$A_L =< W^+ l_\alpha^- | H_{int} | \nu >< \nu | H_{int} | W^- l_\beta^+ > \tag{1.9}$$

The amplitude will be proportional to the elements of the mixing matrix, in the form $U_{\alpha i} U_{\beta i} \lambda_i$ where λ_i are the Majorana phases. However, in this case, the total amplitude is not factorizable. This is explicitly shown in the next section for specific production and detection processes. There occurs a helicity flip from the nature of the neutrino, i.e. as it is its own antiparticle. This helicity flip is proportional to the Majorana mass. We can see this using the Majorana condition $\nu = \lambda\nu^c$, we can rewrite

$$l_{\alpha L}\gamma^{\mu}U_{\alpha i}\nu_{iL} = -\bar{\nu^c}_{iR}\gamma^{\mu}U_{\alpha i}l^c_{\alpha R} = -\bar{\nu}_{iR}\gamma^{\mu}U_{\alpha i}l^c_{\alpha R} \tag{1.10}$$

so that the neutrino field in the left in eq.1.9 can be contracted with the field neutrino field in the right of eq.1.9 in order to make the usual neutrino propagator,$< 0|T(\nu_i\nu_i)|0 >$. Assuming there is no flavor changing the total amplitude will be given by

$$A_{\alpha\alpha} = \left(\sum_i \lambda_i U^2_{\alpha i}\mathrm{m}_i\right) K \tag{1.11}$$

where m_i is the Majorana mass and K is a kinematical and nuclear factor that is not in principle factorizable. The term $(\sum_i \lambda_i U^2_{\alpha i}\mathrm{m}_i)$ is the Majorana effective mass $< m_{\alpha\alpha} >$. Neutrino oscillation experiments can not provide information on the nature of the neutrino, and neutrino antineutrino oscillation experiments have been regarded as unrealistic because of its helicity supression factor m/E. As stated before the Majorana nature of massive neutrinos manifests in lepton number violating processes. The feasible experiments having the potential to establish the Majorana nature of the neutrinos are the neutrinoless double beta decay $(\beta\beta)_{0\nu}$. The observation of this decay and the measurement of its half-life with good precision will not only prove the Majorana nature but will provide information on the neutrino mass spectrum. However measurements of the effective electron-neutrino mass in the neutrinoless double beta decay ($0\nu\beta\beta$) experiments can not restrict the two Majorana CP violating phases present in the PMNS mixing matrix, and neutrino oscillation experiments do not depend on the Majorana phases. This will be a good reason to look for neutrino antineutrino oscillations, were CP violating phases are manifested.

1.2.2 Neutrino antineutrino oscillations

Let us start by considering a positive charged pion which decays into a virtual neutrino at the space-time location (x, t) together with a positive charged muon. After propagating, the neutrino can be converted into an antineutrino which produces a positive charged muon at the point $(x^{'}, t^{'})$ when it interacts with a target, as shown in fig.1.2. For the case of a heavy virtual neutrino one would expect small propagation distances; conversely, long baseline experiments would be sensitive to light neutrinos. In any case, the non-observation of neutrino-antineutrino conversion can be used to set bounds on the effective Majorana mass of the muon neutrino.

For definiteness, we illustrate this process with the production of the neutrino in π^+ decay and its later detection via its weak interaction with a target nucleon N

$$\begin{aligned}\pi^+(p_1) \rightarrow \quad & \mu^+(p_2) \quad +\nu_\mu^s(p) \\ & \hookrightarrow \quad \overline{\nu_l}^d(p) + N(p_N) \rightarrow N'(p_{N'}) + \mu^+(p_l)\end{aligned}$$

where the superscript $s(d)$ refers to the virtual neutrino (antineutrino) at the source (detection) vertex. This $|\Delta L| = 2$ process is a clear signal for neutrino-antineutrino oscillations of the muon type and its amplitude should be proportional to neutrino Majorana masses. It is remarkable that the two identical anti-muons ($\mu^+(p_2)$ and $\mu^+(p_l)$) are produced at very different space-time locations, well separated in distance, thus, there is not any chance to be confused as identical fermions, and the total amplitude does not require to be anti-symmetrized.

If one ignores other flavors, the time evolution of the $\nu_\mu - \bar{\nu}_\mu$ system would be analogous to that of the $K^0 - \bar{K}^0$ or $B^0 - \bar{B}^0$ systems. Instead, we prefer to use the formalism developed in Ref.[27], where the whole reaction includes the production and detection processes of neutrinos. The decay amplitude becomes (for simplicity we assume that leptonic flavor is

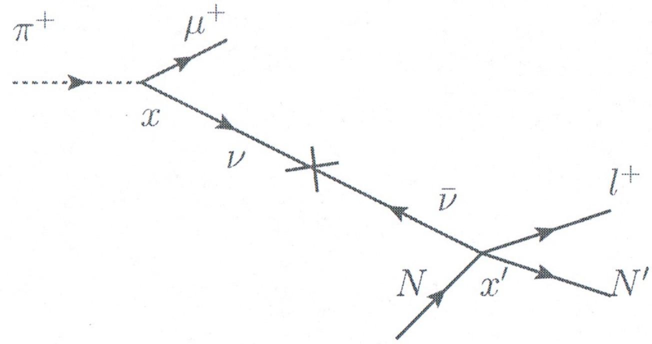

Figure 1.2: Feynman diagram of the process $\pi(p_1) \to \mu^+(p_2)+\nu_\mu(p)$ followed by the detection process:$\overline{\nu_l}(p) + N(p_N) \to N'(p_{N'}) + \mu^+(p_l)$.

conserved at the production and detection vertices):

$$\begin{aligned} T_{\nu_\mu - \bar{\nu}_\mu}(t) &= (2\pi)^4 \delta^4(p_l - p_N + p_{N'} + p_2 - p_1) \\ &\times (G_F V_{ud})^2 (J_{NN'})_\mu f_\pi \\ &\times \sum_i \bar{v}_\mu(p_l)\gamma^\mu(1+\gamma_5)\ \not{p}_1 v(p_2) \\ &\times U_{\mu i} U_{\mu i}(m_{\nu_i}) \frac{e^{-itE_{\nu_i}}}{2E_{\nu_i}}\ , \end{aligned} \tag{1.12}$$

where the relation $\nu_k = \sum U_{k\alpha}\nu_\alpha$ between flavor k and mass α neutrino eigenstates has been used, $f_\pi = 130.4$ MeV is the $\pi^\pm$ decay constant, and $J_{NN'}$ parametrizes the interaction with the nucleon. Note that, contrary to the case of neutrino-neutrino oscillations [27], only the neutrino mass term survives in this case. The only nonvanishing matrix elements of the nucleon are written

$$\begin{aligned} < N(p')|V_\mu|N(p) > &= \bar{u}(p')(\gamma_\mu F_V(q^2) + i\frac{1}{2m_N}\sigma_{\mu\nu}q^\nu F_W(q^2))u(p) \\ < N(p')|A_\mu|N(p) > &= \bar{u}(p')(\gamma_\mu\gamma_5 F_A(q^2))u(p) \end{aligned} \tag{1.13}$$

where $q = p_{N'} - p_N$. When the momentum transfer between the two nucleons is small, as in our case, we may drop the weak magnetism (tensor) F_W and approximate the vector and axial form factors with two constants: $F_V(0) = g_V$ and $F_A(0) = g_A$. The nucleon current then reads,

$$(J_{NN'})_\mu = \overline{u}_{N'}(p_{N'})\gamma_\mu[g_V(q^2) + g_A(q^2)\gamma_5]u_N(p_N) \tag{1.14}$$

If we neglect terms of $O(m_\mu/m_{N,N'})$, one obtains

$$\begin{aligned} \left|T_{\nu_\mu - \bar{\nu}_\mu}(t)\right|^2 &= (2\pi)^4\delta^4(p_l - p_N + p_{N'} + p_2 - p_1)(G_F V_{ud})^4 \\ &\times \; |f_\pi|^2 \sum_{i,j} U_{\mu i}U^*_{\mu j}U_{\mu i}U^*_{\mu j}e^{-it\triangle E_{\nu_{ij}}}\frac{m_{\nu_i}m_{\nu_j}}{4E_{\nu_i}E_{\nu_j}} \\ &\times \; 64(g_A - 1)^2 m_N m_\mu^2(E_2 - E_p) \\ &\times \; \left(\left[1 - \frac{m_N}{(E_2 - E_p)}G(g_A)\right] p_l \cdot p_2 \right. \\ &- \; 2m_N F(g_A)\left[E_2 - E_l\left(1 + \frac{m_\pi^2}{m_\mu^2}\right)\right] \\ &- \left. \frac{1}{2}(m_\mu^2 - m_\pi^2)\right) \end{aligned} \tag{1.15}$$

where $E_2(E_l)$, E_p are, respectively, the initial (final) muon and the pion energies and $\triangle E_{\nu_{ij}} = E_{\nu_i} - E_{\nu_j}$ in the laboratory frame. The functions $F(g_A)$ and $G(g_A)$ are given by: $F(g_A) = \frac{g_A^2+1}{(g_A-1)^2}$, $G(g_A) = \frac{g_A+1}{g_A-1}$. One can easily check that Eq. (1.15) is not factorizable into (production)×(propagation)×(detection) subprocesses due to the terms proportional to $p_l \cdot p_2 = E_l E_2 - |\vec{p_l}||\vec{p_2}|\cos\alpha$, where α is the angle between the directions of μ^+ particles. This is an important difference with respect to the case of neutrino-neutrino ($\Delta L = 0$) oscillations where it was shown in Ref.[27] that the S-matrix formalism reproduces the hypothesis of factorization of the probabilities. The square of the amplitude in terms of invariant quantities is given in appendix A.1.1, as well as a numerical analysis of the kinematical allowed region. The process is seen as a $2 \to 3$ cross section for macroscopically separated particles (which can be neverthelles described in a quantum field theory), this is the $\Delta L = 2$ process $\pi^+ p \to n\mu^+\mu^+$. This phase space can not be solved analitically in closed form, therefore we attemp to make reasonable assumptions for the kinematic allowed regions in the this analysis and attempt to make an estimate on the total cross section in order to obtain a bound over the Majorana mass.

In the following, we shall neglect the q^2-dependence of the nucleon form factors (namely, we take $g_V = g_V(q^2 = 0) = 1$ and $g_A = g_A(q^2 = 0) \approx -1.27$ [7]). As is well known [34], the cross section of charged current neutrino-nucleon quasielastic scattering is sensitive to the q^2-dependence of these form factors. However, as long as we confine to the CP rate asymmetry for neutrino$\leftrightarrow$antineutrino oscillations we expect that the effects of the momentum-transfer dependence of $g_{V,A}$ will partially cancel in the ratio of oscillation rates. Thus, after integration over kinematical variables, it is possible to write the rate of the complete process as

$$\Gamma_{\nu_\mu - \bar{\nu}_\mu} = \left| \sum_i U_{\mu i}^2 \frac{m_{\nu_i}}{2E_{\nu_i}} e^{itE_{\nu_i}} \right|^2 \times F(M, \phi), \tag{1.16}$$

where $F(M, \phi)$ denotes the kinematical function

$$\begin{aligned} F(M,\phi) &= \frac{\pi}{2E_p} (G_F V_{ud})^4 \left| f_\pi \right|^2 64 (g_A - 1)^2 \\ &\times \Big(\Big[I_4 - m_N G(g_A) I_5 - \frac{1}{2}(m_\mu^2 - m_\pi^2) I_1 \Big) \Big] m_\mu^2 \\ &- 2 m_N F(g_A) \left[m_\mu^2 I_2 - (m_\mu^2 + m_\pi^2) I_3 \right] \Big) . \end{aligned} \tag{1.17}$$

The functions I_a for $a = 1, .., 4$ can be obtained from the following integral:

$$I_a = \int \frac{d^3 p_2}{2E_2} \frac{d^3 p_l}{2E_l} (E_2 - E_p) f_a \delta(E_p + E_N - E_{N'} - E_l - E_2), \tag{1.18}$$

with $f_1 = 1$, $f_2 = E_2$, $f_3 = E_l$, and $f_4 = (p_l \cdot p_2)$ and $f_5 = (p_l \cdot p_2)/(E_2 - E_p)$.

There are two interesting limits for this process. At very short times, which means that the detection is very close to the production vertex (short-baseline neutrino experiment), one has, assuming that the $E_{\nu_i} \simeq E_\nu$, that

$$\Gamma_{\nu_\mu - \bar{\nu}_\mu} \simeq \frac{|\langle m_{\mu\mu} \rangle|^2}{E_\nu^2} \times F(M, \phi) \tag{1.19}$$

where $\langle m_{\mu\mu} \rangle$ is the effective Majorana mass for the muon neutrino. In the long time limit which corresponds to a long-baseline neutrino experiment, the oscillation terms cannot be neglected and this process depends on a new combination of phases, mixing angles and masses which could give us complementary information on the neutrinoless double beta decays or on any process that depends exclusively on the effective Majorana mass of the neutrinos. Using this expression for the rate it is possible to get the CP asymmetry (defined in A.1.2) which will depend explicitly on Majorana phases.

$$a_{CP} = \frac{\Gamma_{\nu_\mu - \bar{\nu}_\mu} - \overline{\Gamma}_{\nu_\mu - \bar{\nu}_\mu}}{\Gamma_{\nu_\mu - \bar{\nu}_\mu} + \overline{\Gamma}_{\nu_\mu - \bar{\nu}_\mu}} \tag{1.20}$$

$$= \frac{\sum_{i>j} Im\left(U_{\mu i} U_{\mu i} U^*_{\mu j} U^*_{\mu j}\right) m_{\nu_i} m_{\nu_j} \sin\gamma}{\sum_{i>j} \mathrm{Re}\left(U_{\mu i} U_{\mu i} U^*_{\mu j} U^*_{\mu j}\right) m_{\nu_i} m_{\nu_j} \cos\gamma} \tag{1.21}$$

where $\gamma = \frac{1.27 \Delta m^2_{23} L(km)}{E_\nu (GeV)}$. Here Δm^2_{23} is the difference in the squares of second and third eigenstate neutrino masses, $|\Delta m^2_{23}| = (2.43 \pm 0.13) \times 10^{-3}$ eV^2, and L is the distance between production and detection vertices. Finally, E_ν is the energy of the neutrino beam. It is worth mentioning that the time evolution amplitude for the CP-conjugate process corresponds to the observation of μ^- at the source and at the detector. Therefore, the associated nucleon weak vertex is given by $(J_{N'N})_\mu = \overline{u}_N(p_N)\gamma_\mu(g_V + g_A\gamma_5)u_{N'}(p_{N'})$. In estimating the CP asymmetry in Eq.(1.21), we have assumed that $J_{N'N} \simeq J_{NN'}$.

In the limit of $\theta_{13} = 0$, the Majorana phases $\alpha_{1,2}$ are the only sources of CP violation and hence $Im(U_{\mu i} U_{\mu i} U^*_{\mu j} U^*_{\mu j}) \propto \sin(\alpha_i - \alpha_j)$. For $i = 2$ and $j = 3$ one finds

$$a_{CP} \simeq \tan\left[2(\alpha_2 - \alpha_3)\right] \sin\gamma. \tag{1.22}$$

Thus, in the case of long-baseline neutrino experiment like MINOS where the distance L is given by $L = 735$ km and the energy E_ν is typically around $2 - 3$ GeV [35, 36], one finds that $\sin\gamma \sim \mathcal{O}(1)$. Thus, measuring CP asymmetry will be unavoidable indication for large CP violating Majorana phases.

1.2.3 Experimental constraints on neutrino-antineutrino oscillations

The last two decades have witnessed several experiments that investigate the neutrino-antineutrino transitions. It started in 1982 when the BEBC bubble chamber in the CERN SPS neutrino beam set a limit on $\nu_\mu \to \bar{\nu}_e$ and $\nu_e \to \bar{\nu}_e$ through the search for $\bar{\nu}_e$ appearance. Recently, MINOS [36] has measured the spectrum of ν_μ events which are missing after travelling 735 km. It is these missing events which are the potential source of $\bar{\nu}_\mu$ appearance. In their preliminary analysis, they were able to put a limit on the fraction of muon neutrinos transition to muon anti-neutrinos [26]:

$$P(\nu_\mu \to \bar{\nu}_\mu) < 0.026 \; (90\% \; confidencelevel(c.l.)). \tag{1.23}$$

Assuming CPT, this limit can be written as

$$\frac{\Gamma_{\nu_\mu - \bar{\nu}_\mu}}{\Gamma_{\nu_\mu - \nu_\mu}} \quad < \quad 0.026 \; . \tag{1.24}$$

Using our expression for $\Gamma_{\nu_\mu - \bar{\nu}_\mu}$, and the corresponding rate for neutrino oscillations [27], one gets

$$\left| \sum_i U^2_{\mu i} \frac{m_{\nu_i}}{E_{\nu_i}} e^{itE_{\nu_i}} \right|^2 \quad \lesssim \quad 0.001 \; . \tag{1.25}$$

In the limit of ultra relativistic neutrinos, $E_{\nu_i} \simeq E_\nu (1 + m^2_{\nu_i}/2E^2_\nu)$, and keeping the leading terms in the m_{ν_i}/E_ν terms, we get

$$\left| \sum_i U^2_{\mu i} m_{\nu_i} e^{it\frac{m^2_{\nu_i}}{2E_\nu}} \right|^2 \lesssim 0.001 \times E^2_\nu \; . \tag{1.26}$$

To illustrate the usefulness of this relation, let us consider the general case of 3 generations. In this case, one finds

$$\begin{aligned} 0.001 \times E_\nu^2 &\gtrsim |\langle m_{\mu\mu}\rangle|^2 \\ &- 4\sum_{i>j} \mathrm{Re}\left(U_{\mu i}^2 U_{\mu j}^{*2}\right) m_{\nu_i} m_{\nu_j} \sin^2\frac{\Delta m_{ij} L}{4E_\nu} \\ &- 2\sum_{i>j} \mathrm{Im}\left(U_{\mu i}^2 U_{\mu j}^{*2}\right) m_{\nu_i} m_{\nu_j} \sin\frac{\Delta m_{ij} L}{2E_\nu} \end{aligned} \tag{1.27}$$

Assuming that the only phases that appear in the neutrino mixing matrix are the Majorana phases, it is possible to get a bound on the effective muon-neutrino Majorana mass, only depending on the values of the Majorana phases as the oscillation terms cannot be neglected. In such a case, Eq.(1.27) can be written as:

$$\begin{aligned} 0.001 \times E_\nu^2 &\gtrsim |\langle m_{\mu\mu}\rangle|^2 \pm 4\ \sin\left(\frac{\gamma_{32}}{2}\right) m_{\nu_2} m_{\nu_3} \\ &\times |U_{\mu 2}^2 U_{\mu 3}^{*2}| \sin\left(2\alpha_2 - 2\alpha_3 \pm \frac{\gamma_{32}}{2}\right) \\ &- 4\sin\left(\frac{\gamma_{21}}{2}\right) m_{\nu_2} m_{\nu_1} |U_{\mu 1}^2||U_{\mu 2}^{*2}| \\ &\times \sin\left(\alpha_1 + \frac{\gamma_{21}}{2}\right) \\ &\pm 4\sin\left(\frac{\gamma_{31}}{2}\right) m_{\nu_1} m_{\nu_3} |U_{\mu 1}^2||U_{\mu 3}^{*2}| \\ &\times \sin\left(\alpha_2 \pm \frac{\gamma_{31}}{2}\right) \end{aligned} \tag{1.28}$$

where $\gamma_{ij} = \frac{\Delta m_{ij} L(km)}{2E_\nu(GeV)}$, and the positive and negative signs refer to normal and inverted hierarchies, respectively.

We can further neglect $\sin\gamma_{21} \approx 0$ such that $m_{\nu_1} \approx m_{\nu_2}$ at first order in $\mathcal{O}(\frac{\Delta m_{12}^2}{4E_\nu})$ for the fixed experimental parameters in MINOS, then Eq.(1.28) can be written as

$$\begin{aligned} 0.001 \times E_\nu^2 &\gtrsim |\langle m_{\mu\mu}\rangle|^2 + A(\alpha_1, \alpha_2) m_{\nu_2} m_{\nu_3} \\ &\gtrsim |\langle m_{\mu\mu}\rangle|^2 \end{aligned} \tag{1.29}$$

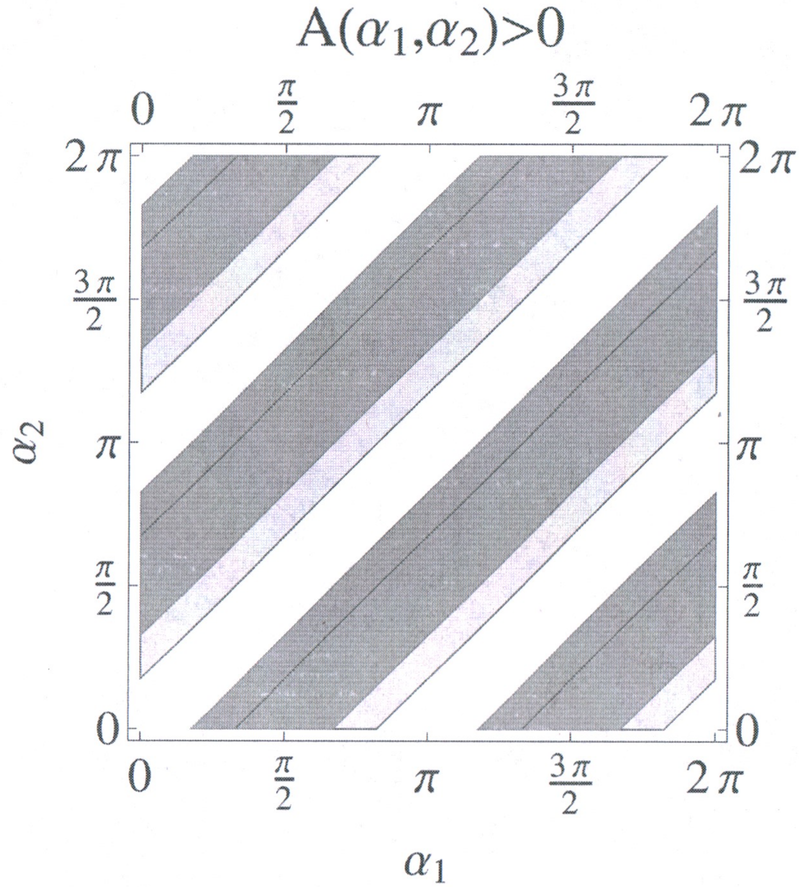

Figure 1.3: Region of the parametric space (α,β) for which the $A(\alpha, \beta)$ coefficient is positive. The light and dark gray zones correspond to an inverted and normal hierarchy schemes respectively.

where the coefficient $A(\alpha_1, \alpha_2)$ is a function of the Majorana phases. In Fig. (1.3) we show the regions in which $A(\alpha_1, \alpha_2) > 0$ is satisfied for both cases of normal or inverted hierarchies, and therefore we can find a stringent bound on the effective Majorana mass.

Thus, using $E_\nu \approx 2$ GeV, one gets the following bound on

$$|\langle m_{\mu\mu}\rangle| \lesssim 64 \text{ MeV}$$

Over the excluded regions it is not possible to get a conservative bound on $|\langle m_{\mu\mu}\rangle|$ without making extra assumptions on the neutrino mass matrix, however Eq.(1.27) can be used to bound Majorana parameters (masses and phases) which appear in $|\langle m_{\mu\mu}\rangle|$. As an example, if $0 \leq 2(\alpha_2 - \alpha_3) \leq \pi - \gamma/2$ and assuming that the effective muon-neutrino Majorana mass is dominated by m_{ν_2} and m_{ν_3} (the two flavour limit case) , it is still possible to get the following conservative bound:

$$|\langle m_{\mu\mu}\rangle| \lesssim 109\ MeV$$

So, within the hypothesis done on Majorana phases, the limits obtained improve by various order of magnitude the present bounds on $\langle m_{\mu\mu}\rangle$ coming from direct searches in $K^+ \to \pi^-\mu^+\mu^+$ decay: $|\langle m_{\mu\mu}\rangle| \leq 0.04$ TeV [15]. Clearly, these bounds lie far above the limits expected from indirect bounds obtained from oscillation experiments and nuclear beta and neutrinoless double beta decays. However, this is the first "direct" experimental limit on the effective mass of muon Majorana neutrinos. The limit might be useful for constraining large non-standard contributions. Also it is expected that some better bounds can be obtained in the future from improved experimental analysis and with more accumulated data sets. Note also that the bound obtained above for $|\langle m_{\mu\mu}\rangle|$ is a factor of 3 above the trivial kinematical bound $m_\pi - m_\mu \approx 34$ MeV that is allowed for the (on-shell) muon neutrino in pion decay. However, this kinematical bound applies only to the effective mass of a lepton-number conserving muon neutrino ($\pi^+ \to \mu^+\nu_\mu$).

In MINOS and, in general, in all long-baseline neutrino experiments, the oscillation terms are not negligible. So, it means that such analysis could not only give us information on the neutrino effective Majorana mass but it could be used to determine parameters of the mixing matrices and get bounds on the absolute value of Majorana masses. Also, if in a long-baseline neutrino experiment neutrino detectors are located at different distances from the source, it should be possible to get enough constraints on the mixing parameters and Majorana masses to fix them.

A bound on the effective Majorana mass of the muon neutrino, which is independent of the mass hierarchies and Majorana phases, can be obtained using the fluxes of ν_μ and $\bar{\nu}_\mu$ measured with the near detector of the MINOS experiment [37]. Since the near detector is located L=1.04 km away from the target and for neutrino energies above 1 GeV, all oscillatory terms in Eq. (5) are equal to 1. Under the assumption that the excess of $\bar{\nu}_\mu$ events arises from $\nu_\mu \to \bar{\nu}_\mu$ transitions we get (note the muon-neutrino and muon-antineutrino total cross sections induced by charged currents are flat for neutrino energies above 2 GeV [7]):

$$\frac{|\langle m_{\mu\mu}\rangle|^2 \int \frac{dE_\nu}{E_\nu^2}}{\int dE_\nu} \leq \frac{\int (\Phi_{\bar{\nu}_\mu}^{Obs}(E_\nu) - \Phi_{\bar{\nu}_\mu}^{MC}(E_\nu))dE_\nu}{\int \Phi_{\nu_\mu}^{Obs}(E_\nu)dE_\nu} . \tag{1.30}$$

where $\Phi_{\nu,\bar{\nu}}^{Obs(MC)}$ denote the observed(expected) fluxes. Using the expected and measured integrated fluxes by the MINOS collaboration for the energy region $5 \leq E_\nu \leq 50$ GeV, we get the following bound:

$$|\langle m_{\mu\mu}\rangle| \leq 2.7\,\mathrm{GeV}, \tag{1.31}$$

which looks much less restrictive than the value reported above.

1.3 New Physics in $\Delta L = 2$ neutrino-antineutrino oscillations

As an illustrative exercise we consider next a virtual neutrino(antineutrino) produced together with a positively(negatively) charged muon at the space time location (x, t), it travels to (x', t') and is detected there because it interacts with a target producing a positively(negatively) charged lepton. We are assuming this $\Delta L = 2$ process is due to NSI interactions at the production vertex. For definitness, we illustrate this process with production of an antineutrino(neutrino) in the $\pi^+(\pi^-)$ decay and its late detection via its weak interaction with a target nucleon N.

$$\pi^+(p_1) \to \mu^+(p_2) + \nu_s^c(p) \hookrightarrow \nu_d^c(p) + N(p_N) \to N'(p_{N'}) + l^+(p_l) \tag{1.32}$$

These effective states are not necessarily ΔL conserving once NSI interactions are introduced. If we assume that neutrinos are left-handed, $\Delta L = 2$ semileptonic interactions can be described by the following effective Hamiltonian,

$$\begin{aligned} \mathcal{H} &= 2\sqrt{2} G_F V_{ud} \left\{ C_1^k \left(\overline{\nu_k^c} \gamma^\alpha P_R \mu\right) \left(\bar{d} \gamma_\alpha P_R u\right) + C_2^k \left(\overline{\nu_k^c} \gamma^\alpha P_R \mu\right) \left(\bar{d} \gamma_\alpha P_L u\right) \right. \\ &\quad \left. + C_{(3,4)}^k \left(\overline{\nu_k^c} P_L \mu\right) \left(\bar{d} P_{(R,L)} u\right) + C_{5(R,L)}^k \left(\overline{\nu_k^c} \sigma_{\alpha\beta} P_R \mu\right) \left(\bar{d} \sigma^{\alpha\beta} P_{(R,L)} u\right) \right\} , \end{aligned} \tag{1.33}$$

where ν_k denotes a neutrino with flavor k and $P_{R,L} = (1 \pm \gamma_5)/2$.

In the following and for simplicity, we consider the case where lepton number violation occurs only at the π^+ decay vertex. Note that the tensor currents proportional to the $C_{5L(R)}^k$ Wilson coefficients will not contribute to π^+ decay because it is not possible to generate an antisymmetric tensor from the pion momentum alone. Thus, the only non-vanishing hadronic matrix elements at the production vertex are:

$$\langle 0 | \bar{d} \gamma^\mu \gamma_5 u | \pi^+ \rangle = i f_\pi p_\pi^\mu \ , \ \langle 0 | \bar{d} \gamma_5 u | \pi^+ \rangle = \frac{-i f_\pi m_\pi^2}{m_u + m_d} \ , \tag{1.34}$$

where $f_\pi = 130$ MeV is the pion decay constant and $m_{u,d}$ denote the light quark masses. The time evolution amplitude for the corresponding process is given by:

$$\begin{aligned} T_{\bar{\nu}_\mu - \bar{\nu}_l}(\tau) &= (2\pi)^4 \delta^4(P_F + p_2 - p_1)(G_F V_{ud})^2 (J_{NN'})_\mu \\ &\quad \times f_\pi \sum_k \overline{u}_l(p_l) \gamma^\mu (1 - \gamma_5)(\not{p}_N - \not{p}_l - \not{p}_{N'}) \\ &\quad \times (m_\pi A_k + B_k \not{p}_\pi) v(p_2) \\ &\quad \times \sum_i U_{li}^* U_{ki} \frac{e^{-i\tau(E_{\nu_i})}}{2E_{\nu_i}} \ . \end{aligned} \tag{1.35}$$

Again we shall neglect the q^2-dependence of the nucleon form factors (namely, we take $g_V = g_V(q^2 = 0) = 1$ and $g_A = g_A(q^2 = 0) \approx -1.27$ [7]). In this case the total rate of the $\Delta L = 2$ is factorized:

$$|T_{\nu_\mu - \nu_\mu^c}(\tau)|^2 \simeq \Gamma_\pi \frac{1}{2}(A^2 m_\pi^2 + B^2 m_\mu^2 - 2\Re[AB^*]) P(\nu_\mu^c \to \nu_\mu^c)\sigma_\nu$$

with $\tau = (t' - t) > 0$ is the time elapsed from the production to the detection space-time locations of neutrinos. The A, B coefficients are given by,

$$B_\mu \equiv -i(C_1^{\mu *} - C_2^{\mu *}) \ , \ A_\mu = \frac{-im_\pi^2}{m_u + m_d}(C_3^{\mu *} - C_4^{\mu *}) \tag{1.36}$$

We shall attempt to constrain the LNV parameters with the MINOS preliminary results [26]: $P(\nu_\mu \to \bar{\nu}_\mu) < 0.026 \sim (90\%\ c.l.)$., this is,

$$(A^2 m_\pi^2 + B^2 m_\mu^2 - 2\Re[AB^*]) \lesssim 0.05 \tag{1.37}$$

As an example let us apply the previous formalism to the MSSM without R-Parity conservation, but requiring the conservation of baryon number to ensure the proton is stable. R-Parity is a discrete symmetry defined as $(-1)^{3B+L+2S}$, where B, L and S are the baryon number, lepton number and particle spin respectively. R-Parity violating (RPV) interactions involve either lepton number violation or baryon number violation, but not both in order to preserve proton stability. These interactions lead to flavor violating interactions in the leptonic and hadronic sector. A vast majority of observables have been used to set the corresponding bounds to these effective couplings (for a complete review see [29] and references therein). The superpotential that conserves bayon number is

$$W_{\Delta L=1} = \lambda_{ijk} L_i \cdot L_j \bar{e}_k + \lambda'_{ijk} L_i \cdot Q_j \bar{d}_k + \mu_L^i L_i \cdot H_u \tag{1.38}$$

here the "$\cdot$" is a $SU(2)$ product, therefore λ_{ijk} is antisymmetric in i, j. The summation over the generation indices $i, j, k = 1, 2, 3$ is understood. The bar symbol represents an antiparticle

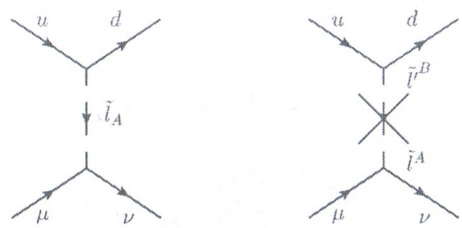

Figure 1.4: $\Delta L = 2$ SUSY contribution to the $\pi^+ \to \mu^+ + \nu^c$

and not the Dirac conjugation. We can expand the Yukawa terms in the Lagrangian as

$$\begin{aligned} L &= \lambda_{ijk}\left[\tilde{\nu^i}_L \overline{e^k_R} e^j_L + \tilde{e^j}_L \overline{e^k_R} \nu^i_L + e^{*k}{}_R \overline{(\nu^i_L)^c} e^j_L - (i \leftrightarrow j)\right] \\ &+ \lambda'_{ijk}\left[\tilde{\nu}^i_L \overline{d^k_R} d^j_L + \tilde{d}^j_L \overline{d^k_R} \nu^i_L + \tilde{d^*}^k_R \overline{(\nu^i_L)^c} d^j_L - \tilde{e}^i_L \overline{d^k_R} u^j_L \right. \\ &- \left. \tilde{u}^j_L \overline{d^k_R} e^i_L - \tilde{d^*}^k_R \overline{(e^i_L)^c} u^i_L\right] + h.c. \end{aligned} \tag{1.39}$$

The corresponding effective hamiltonian (Fig.3) is therefore,

$$H_{eff} = \frac{G_F}{\sqrt{2}} V_{ud} \sum_i \left((\bar{\nu}\gamma^\rho P_L \mu)(\bar{d}\gamma_\rho P_L u) + C_2(\bar{\nu} P_R \mu)(\bar{d} P_L u) + C_3(\overline{\nu^c} P_L \mu)(\bar{d} P_L u)\right) \tag{1.40}$$

The Wilson coefficients can be expressed as $C_i = C_i^{SM} + C_i^{SUSY}$. For $i = 2, 3$ the C_i^{SM} vanish identically, and $C_1 = 1$. In this respect, the Wilson coefficients C_i are given by

$$C_2 = \frac{\sqrt{2}}{G_F V_{ud}} (\frac{1}{\tilde{m}^2})(\lambda^*_{122}\lambda'_{211} + \lambda_{232}\lambda_{311}) \ , \ C_3 = \frac{\sqrt{2}}{G_F V_{ud}} (\frac{1}{\tilde{m}^4}) \sum_{A,B} \lambda_{12A}\lambda'_{B11}(\delta^l_{RL})_{AB} \tag{1.41}$$

Note however that only C_3 describes the $\Delta L = 2$ neutrino conversion (see eq. 1.33).

Using our previous results (Eq. 1.37) and assumptions, for typical slepton masses of 200 GeV, we obtain $\sum_{A,B} \lambda_{12A}\lambda'_{B11}(\delta^l_{RL})_{AB} \lesssim 10^{-4}$. This bound is as expected less restrictive than previous bounds [7].

Chapter 2

New Physics with charmed mesons D

2.1 Introduction

In spite of the Standard Model (SM) success, now favored by the probable recent discovery of the Higgs boson [1, 2], the search of a more fundamental theory at an energy scale much higher than the electroweak scale is still open. But even at lower energies, comparable with the electroweak scale there are still fundamental aspects of the theory that are yet to be answered, such as the size of CP violation, which we know nature exhibits. Interestingly, low energy scale experiments may shed some light in the search for such fundamental theory due to their possibility of getting high statistics and hence indirect observables of New Physics (NP). We will use D meson decays as an illustration. Contrary to B meson physics, charmed hadronic states are in the unique mass range of $O(2\,\mathrm{GeV})$, which allows for strong non perturbative hadronic physics [38]. Moreover, the calculations for the relevant form factors, which parameterize most of the QCD effects within the hadronic state, have been improved significantly reaching an acceptable precision of $\sim$2% [39, 40, 6]. The SM predictions for the D meson leptonic and semileptonic decays relies on the lattice QCD estimates of the form factors, and appear to be in agreement with the world average experimental measurements [6], allowing us to disprove New Physics hypothesis or find restrictive bounds over several models beyond the SM.

At low energies, most of the extensions to the Standard Model reduce to an effective four Fermi interaction, usually called Non Standard Interaction NSI, that can be parameterized by a generic coefficient (Fig. 2.1). For the $\Delta C = \Delta S$ leptonic and semileptonic D meson decays, the new particle state should couple to the leptons and the second generation of quarks, leaving such effective interaction. Any kind of intermediate state, such as scalars, vectors or even tensors, are allowed. Examples are the Two Higgs Doublet Model Type-II (THDM-II) and Type III (THDM-III) [41], the Left-Right model (LR)) [42, 43, 44, 45, 46], the Minimal Supersymmetric Standard Model with explicit R-Parity violation (MSSM-$\not{R}$) [28, 29], and the Leptoquark model [47, 48], also illustrated in Fig. 2.1.

Non Standard interactions from a model independent approach had been considered and constrained with D_s leptonic decays [52, 53], and independently, using semileptonic decays [54, 53]. In this work we make a model independent analysis and a model dependent analysis in order to constrain NSIs combining the leptonic and semileptonic decays of the D meson [5]. We use the latest Lattice results on the form factors[6] which have reached a significant precision. We show the usefulness of the model independent constraints as well as specific cases when a model dependent analysis is needed. The q^2 distributions for the $D^+ \to \bar{K}^0 e^+ \nu_e$ and $D^0 \to K^- e^+ \nu$ decays, which are expected to be sensitive to new physics, are also considered. Using the respective bounds for the Wilson coefficients, we compute as well the transverse polarization of the charged lepton in the semileptonic decay of the D meson. This T violating observable has not been measured but may provide significant constraints over the complex character of the new physics parameters, as in the case of the B meson semileptonic decay [49] and other meson decays [50]. This chapter is organized as follows: In Section 2.3 we describe the general effective Lagrangian for the semileptonic transition $c \to s$ when non standard interactions are included and show the theoretical branching ratios

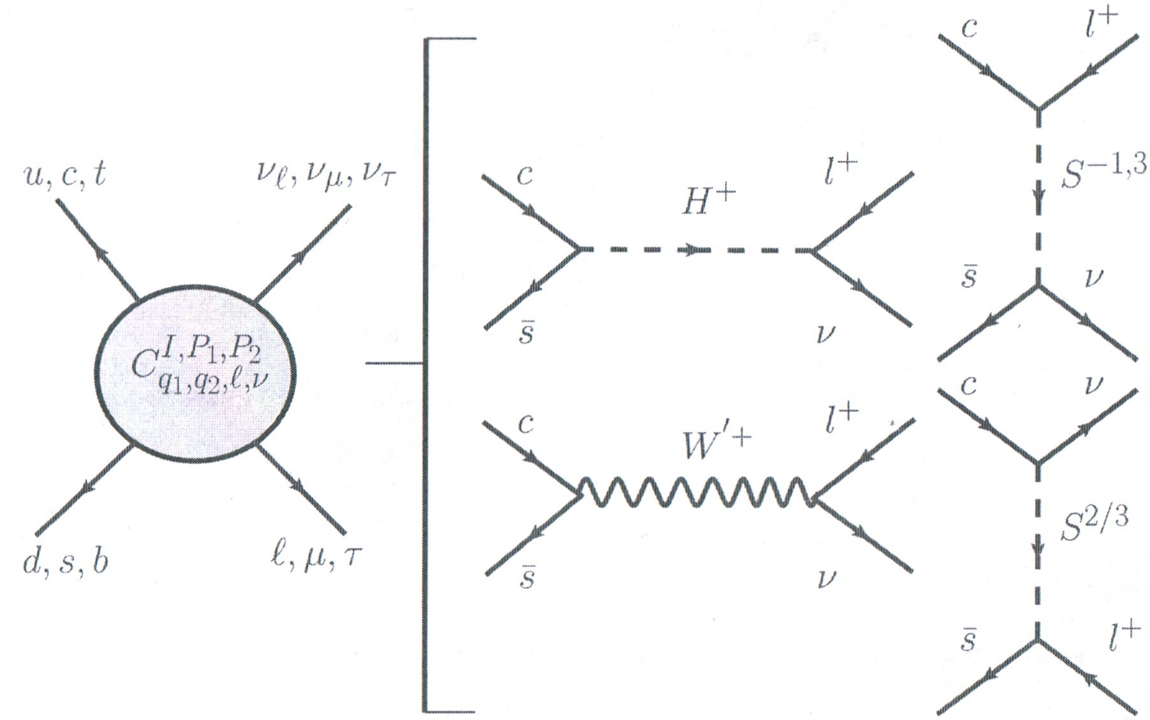

Figure 2.1: Generic charged current non standard interaction between two quarks and the leptonic sector. Some Feynman diagrams for models beyond SM involving the $c\bar{s} \to l\bar{\nu}$ transition involved in D meson decays are shown.

and the transverse polarization of the D meson semileptonic decay. In Section 2.4 we show the experimental constraints over the Wilson coefficients, and the theoretical predictions for the transverse polarization of the D meson semileptonic decay. In Section 2.5 we constrain the relevant parameters of the THDM-II and THDM-III, LR, and the MSSM-$\not{R}$, and the leptoquark model. In Section ?? we give our conclusions and comments over the relevance of these bounds.

2.2 D meson leptonic and semileptonic decays in the SM

We will briefly review the leptonic $D_s \to \ell_i \nu_i$ and semileptonic $D^\pm \to K^0 \ell_i^\pm \nu_i$, $D^0 \to K^\pm \ell_i^\mp nu_i$ decays of the D meson in the SM. The D meson is a pseudoscalar composite particle of two quark states, such as $D_s^+ = c\bar{s}$, $D^0 = c\bar{u}$, $D^+ = c\bar{d}$ and their corresponding antiparticles. The quarks however interact strongly; such interactions are well described by the QCD (Quantum Chromodynamics) gauge symmetry and assigns an extra quantum number to the quarks, 'the color charge'. At low energies, in the range where the D mesons are produced, the strong interaction among the quarks is not perturbative, in fact they may show confinement. One may parameterize our ignorance of the detailed structure of the strong interactions with *form factors*. These form factors can be measured in the experiments, or calculated from first principles within a non-perturbative approach, e.g. Lattice QCD (LQCD) [51], a lattice gauge theory formulated on discrete spacetime. As mentioned in the introduction, we will consider the LQCD results relevant for the D meson form factors, the details of their calculations are beyond the scope of this work.

The quark couplings to the W-boson is given within the Standard Model by the following interaction Lagrangian,

$$L_{int} = -\frac{g}{\sqrt{2}} W_\mu^+ \sum_{ij} \bar{u_i}_L \gamma^\mu V_{ji}^* d_{jL} - \frac{g}{\sqrt{2}} W_\mu^- \sum_{ij} \bar{d_i}_L \gamma^\mu V_{ji} u_{iL} \tag{2.1}$$

Here, g is the semiweak coupling constant, V_{ji} is the CKM matrix element (analogous to the PMNS mixing matrix of neutrinos 1.5) where the indexes i, j run over the quark flavours $i = u, c, t$ and $j = d, s, b$. The mixing matrix, has 3 angles and one CP violating phase, as in the case of Dirac neutrinos. Thus the effective Hamiltonian for the leptonic (two body) decay of the D_s meson is given by

$$H_{eff} = 2\sqrt{2}G_F V_{cs}^*(\bar{s}_L\gamma_\mu c_L)(\bar{\nu}_L\gamma^\mu \ell_L) \tag{2.2}$$

As in the case of the pion, we can only construct an axial current with a single momenta for the pseudoscalar D_s meson (within the SM), therefore the only non-vanishing hadronic matrix element is $< 0|\bar{s}\gamma^\mu\gamma^5 c|D_s(p) >\equiv if_{D_s}p^\mu$. This normalisation is associated with the pion decay constant $f_\pi = 131\,\text{MeV}$. The constant f_{D_s} is the D_s decay constant.

The matrix element of the first order perturbation of the effective Hamiltonian between the initial and final states is given by

$$M_{if} =< \nu_l(p_1)\ell(p_2)|H_{eff}|D_s(p) >= -if_{D_s}\frac{G_F V_{cs}^*}{\sqrt{2}}\bar{u}(p_1)(1+\gamma^5)\not{p}v(p_2) \tag{2.3}$$

The amplitude square is easily calculated to be $|M|^2 = 4|f_{D_s}G_F V_{cs}|^2 m_l p_1 \cdot p_2$. The decay rate is given by

$$d\Gamma = \frac{1}{m_{D_s}}|M|^2(2\pi)^4\delta^4(p-p_1-p_2)\frac{d^3p_1}{(2\pi)^3 2p_1^0})\frac{d^3p_2}{(2\pi)^3 2p_2^0} \tag{2.4}$$

It is straight foreward to integrate over the neutrino momentum vector. Choosing the decay reference system, i.e. $p = (m_{D_s}, 0)$, and assuming approximately massless neutrinos, this leads to

$$\Gamma = \frac{|f_{D_s}G_F V_{cs}^*|^2}{8\pi^2 m_{D_s}}\int \delta(m_{D_s} - E_l - \sqrt{E_l^2 - m_l^2})dE_l d\Omega(\frac{m_{D_s}^2}{2})(1 - \frac{m_l^2}{m_{D_s}^2}) \tag{2.5}$$

Using the identity $\delta(f(x)) = \delta(x - x_0)/|f'(x_0)|$ we have

$$\Gamma = \frac{|f_{D_s}G_F V_{cs}^*|^2}{8\pi}m_l^2 m_{D_s}(1 - \frac{m_l^2}{m_{D_s}^2})^2 \tag{2.6}$$

Now let us consider the decay of the pseudoscalar meson into 3 bodies $D(p)^+ \to K(k)^0\ell(p_2)^+\nu(p_1)$. The hadronic matrix element in this case is written as

$$< K(k)|\bar{s}\gamma^\alpha c|D(p) >= (p^\alpha + k^\alpha - \frac{m_D^2 - m_K^2}{q^2}q^\alpha)f_+(q^2) + \frac{m_D^2 - m_K^2}{q^2}q^\alpha f_0(q^2) \tag{2.7}$$

where $q^2 = (p-k)^2$. In the massless lepton limit the $f_0(q^2)$ contribution vanishes. The amplitude of the decay process is then given by

$$\begin{aligned} A(D(p) \to K(k)\ell(p_2)\nu(p_1) &= \frac{G_F}{\sqrt{2}} V_{cs}^* \bar{u}(p_1)(1+\gamma_5)\left[\not{Q} f_+(q^2)\right. \\ &+ \left.\left(\frac{m_D^2 - m_K^2}{q^2}\right)(f_0(q^2) - f_+(q^2))\not{q}\right] v(p_2) \end{aligned} \tag{2.8}$$

where $Q = p + k$. It follows that the amplitude square in terms of invariant quantities is

$$\begin{aligned} |A_{D\to k\nu l}|^2 &= -4G_F^2|V_{cs}|^2 \left[p_1 \cdot p_2 \left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+)\right.\right. \\ &\times \left.((m_D^2 - m_K^2)(f_0 - f_+) + 2f_+ q \cdot Q) + |f_+|^2 Q^2\right) \\ &- 2\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+)p_2 \cdot q + f_+ p_2 \cdot Q\right) \\ &\times \left.\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+)q \cdot p_1 + f_+ Q \cdot p_1\right)\right] \end{aligned} \tag{2.9}$$

In the rest frame (RF) of the decaying meson, the differential decay rate for example for the $D^0 \to K^\pm l^\mp \nu$ decay channel, after a trivial integration over the neutrino phase space, is given by

$$\begin{aligned} \frac{d\Gamma}{dE_K dE_{l_\perp}} &= |G_F V_{cs}|^2 m_D/(2\pi)^3 \left(((E_K^2 - m_K^2)(1 - m_l^2/q^2) - 4E_{l_\perp}^2)f_+(q^2)^2\right. \\ &+ \frac{(q^2 - m_l^2)}{4m_D^2}\left(\frac{(m_D^2 - Mm_K^2)}{q^2}\right)^2 m_l^2 * f_0(q^2)^2 \\ &- \left. 2\frac{m_l^2}{m_D}\frac{(m_D^2 - m_K^2)}{q^2} E_{l_\perp} f_0(q^2) f_+(q^2)\right) \end{aligned} \tag{2.10}$$

where $q^2 = m_D^2 + m_K^2 - 2m_D E_K$ and $E_{l_\perp} = El - \frac{1}{2}(m_D - E_K)(1 + m_l^2/q^2)$ such that there are only two Lorentz independent invariants $E_l = p \cdot p_2/m_D$ and $E_K = p \cdot k/m_D$ in the RF. The kinematical allowed regions [54] are defined by $-\frac{1}{2}(1 + m_l^2/q^2)\sqrt{E_K^2 - m_K^2} < E_{l_\perp} <$

$\frac{1}{2}(1 - m_l^2/q^2)\sqrt{E_K^2 - m_K^2}$ and $m_K < E_K < (m_D^2 + m_K^2 - M_l^2)/2m_D$. After integrating over the variable $E_{l_\perp}$ we have

$$\begin{aligned} \frac{d\Gamma}{dE_K} &= \frac{G_F^2|V_{cs}|^2}{(2\pi)^3 3m_D}\sqrt{E_K^2 - m_K^2}\left(1 - \frac{m_l}{q^2}\right)^2 \left(m_D^2|f_+(q^2)|^2(E_K^2 - m_K^2)\frac{2q^2 + m_l^2}{q^2}\right. \\ &+ \left.\frac{3}{4}|f_0(q^2)|^2\frac{m_l^2}{q^2}(m_D^2 - m_K^2)^2\right) \end{aligned} \tag{2.11}$$

Where we use the same functional shape of the form factors as in [66]:

$$\begin{aligned} f_+(q^2) &= \frac{f_+(0)}{(1 - \frac{\alpha q^2}{M_X^2})(1 - \frac{q^2}{M_X^2})} \\ f_0(q^2) &= \frac{f_+(0)}{(1 - \frac{q^2}{M_X^2\beta})} \end{aligned} \tag{2.12}$$

where M_X is the appropiate pole mass, i.e. $M_{D_s^*}$ for f_+ and $M_{D_{s0}^*}$ for f_0; $\alpha = 0.5$ and $\beta = 1.31$ [66]. At $q^2 = 0$, $f_+(0) = 0.745 \pm 0.011$ [6].

In fig. 2.2 we show the difference in the differential decay rate when the charged lepton mass is neglected. We can see that as one would expect the electron mass is in a good approximation negligible but not the muon mass. In the following we take into account the complete defferencial decay rate. In appendix A.2.1 we derive a common expression found in literature in the massless limit for a different choice of reference system.

2.3 Non standard interactions and relevant observables

2.3.1 Effective Lagrangian below the electroweak scale

The search of new physics effects in the leptonic and semileptonic processes of mesons has two sources of uncertainty that can not be separated: the non perturbative long-distance forces that bind quarks forming hadrons and the determination of the free parameters of the SM, *i.e.* quark masses and CKM matrix elements. The non-perturbative QCD effects

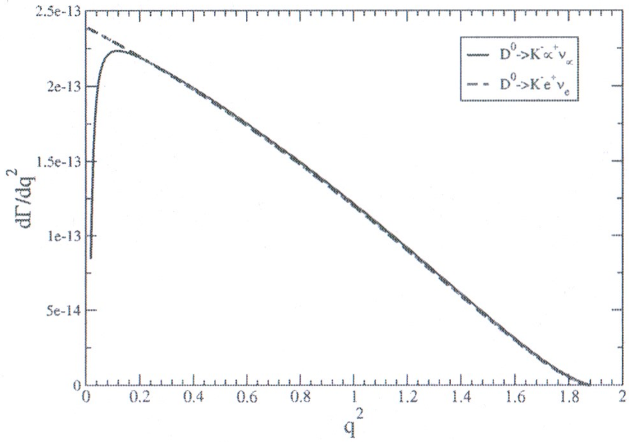

Figure 2.2: Decay rate for $D^0 \to K^- \nu_\ell \ell^+$. The dashed blue line is the diffential decay rate for massless charged leptons.

are parameterized introducing form factors. On the other hand effective couplings that correspond to short distance interactions could receive non-standard contributions. Hence flavour-changing meson transitions in the SM have at least two scales involved, the electroweak scale that is responsible of the flavour changing and the scale of strong interactions [55]. When NSI are considered, we assume that the new physics energy scale is higher than the electroweak scale, thus the operator product expansion formalism (OPE) [56] is suitable since it allows the separation between long-distance (low energy) and short-distance (high energy) interactions. In the OPE the degrees of freedom corresponding to higher energies scales are integrated out [57], resulting an effective Lagrangian where all high energy physics effects are parameterized by Wilson's coefficients, namely the effective couplings multiplying the operators of the Lagrangian. Let us consider then the non-standard effective Lagrangian for a semileptonic transition as the one illustrated in Fig. 2.1 is:

$$-\frac{\mathcal{L}_{NP}}{G_F} = \sum_{\substack{c,s,\ell,\nu \\ I=S,V,T \\ P_{1,2}=L,R}} C^{I,P_1P_2}_{q_1q_2\ell\nu} (\bar{q}_1 \Gamma^I P_1 q_2) \cdot (\bar{\nu}_L \Gamma_I P_2 \ell) \,, \tag{2.13}$$

where the indexes q_1 and q_2 represent down-type and up-type quarks respectively, ℓ is the charged lepton flavor and ν its corresponding neutrino. $P_{1,2}$ represent the chiral projectors $L = (1-\gamma^5)/2$ and $R = (1+\gamma^5)/2$. Here, the current operators Γ's are determined by the Dirac field bilinears, namely: $\Gamma_S = 1$, $\Gamma_V = \gamma_\mu$ and $\Gamma_T = (i/2)[\gamma^\mu, \gamma^\nu]$. The dimensionless coefficients $C^{I,P_1P_2}_{q_1q_2\ell\nu}$ have a clean interpretation: they are a measurement of how big can the NSI be as compared to the SM current, since they are weighted by the Fermi constant G_F. This parametrization technique enables us to test NSI when the experiments reach certain precision, and in particular to look for NP effects at low energies. In particular, for the $c\bar{s} \to \bar{\nu}\ell$ transition, and considering only left handed neutrinos, the non-standard effective Lagragian reads,

$$\begin{aligned}
-\frac{\mathcal{L}_{NP}}{G_F} &= C^{V,LL}_{q_1q_2\ell\nu}(\bar{s}_L\gamma_\mu c_L)(\bar{\nu}_L\gamma^\mu l_L) \\
&+ C^{V,RL}_{sc\ell\nu}(\bar{s}_R\gamma_\mu c_R)(\bar{\nu}_L\gamma^\mu l_L) \\
&+ C^{S,RR}_{sc\ell\nu}(\bar{s}_L c_R)(\bar{\nu}_L l_R) \\
&+ C^{S,LR}_{sc\ell\nu}(\bar{s}_R c_L)(\bar{\nu}_L l_R) \\
&+ C^{T,LR}_{sc\ell\nu}(\bar{s}_R\sigma_{\mu\nu} c_L)(\bar{\nu}_L\sigma^{\mu\nu} l_R) \\
&+ C^{T,RR}_{sc\ell\nu}(\bar{s}_L\sigma_{\mu\nu} c_R)(\bar{\nu}_L\sigma^{\mu\nu} l_R)
\end{aligned} \tag{2.14}$$

2.3.2 NSIs in the D meson leptonic and semileptonic decays

First, let us analyse the full leptonic decay $D_s \to \ell\nu_\ell$ with effective non-standard interactions. We can not construct an antisymmetric tensor with a single momenta, therefore the only non-vanishing hadronic elements for a pseudoscalar meson decay into two leptons are

$$< 0|\bar{s}\gamma_\mu\gamma_5 c|D(p) >= i f_{D_s} p_\mu \tag{2.15}$$

$$< 0|\bar{s}\gamma_5 c|D(p) >= i f_{D_s}\frac{m_{D_s}}{m_c + m_s} \tag{2.16}$$

Then we get the following amplitude

$$< \bar{\nu}_l(p_1)l(p_2)|\mathcal{H}_{NP}|D_s(p) >= \bar{u}(p_1)\mathcal{O}v(p_2) \tag{2.17}$$

where the operator $\mathcal{O}$ is given by

$$\mathcal{O} = \frac{f_{D_s}}{\sqrt{2}}(1+\gamma_5)G_F\left(i\left[\frac{(C^{V,LL}_{sc\ell\nu} - C^{V,RL}_{sc\ell\nu})}{2\sqrt{2}} - V^*_{cs}\right]\not{p}_{D_s} + i\frac{m^2_{D_s}(C^{S,RR}_{sc\ell\nu} - C^{S,LR}_{sc\ell\nu})}{2\sqrt{2}(m_c+m_s)}\right) \tag{2.18}$$

It follows that the amplitude square is

$$|A(D\to \ell\nu|^2 = 4G_F^2|V^*_{cs}m_l + m_l\frac{(C^{V,LL}_{sc\ell\nu} - C^{V,RL}_{sc\ell\nu})}{2\sqrt{2}} + m_{D_s}\frac{(C^{S,RR}_{sc\ell\nu} - C^{S,LR}_{sc\ell\nu})}{2\sqrt{2}(m_c+m_s)}|^2(p_2\cdot p_1) \tag{2.19}$$

Thus, in the rest frame of the decaying meson the decay rate of $D_s \to \ell\nu_\ell$ including the SM Lagrangian plus the NSI Lagrangian of eq. (2.12), is given by

$$\Gamma_{D_s\to\ell\nu} = \frac{|G_F f_{D_s}\left(M^2_{D_s} - m_l^2\right)|^2}{8\pi M^3_{D_s}}\left|V_{cs}m_l + \frac{m_l(C^{V,LL}_{sc\ell\nu} - C^{V,RL}_{sc\ell\nu})}{2\sqrt{2}} + \frac{M^2_{D_s}(C^{S,RR}_{sc\ell\nu} - C^{S,LR}_{sc\ell\nu})}{2\sqrt{2}(m_c+m_s)}\right|^2.$$

For the semileptonic decay of the D meson, e.g. $D(p)^0 \to K^{\mp}\nu(p_1)l(p_2)^{\pm}$, where we have two definite four-momenta, we will have a contribution coming from the tensor current too. The non-vanishing hadronic elements are

$$\begin{aligned} < K(k)|\bar{s}\gamma_\mu c|D(p) > &= (p^\alpha + k^\alpha - \frac{m_D^2 - m_K^2}{q^2}q^\alpha)f_+(q^2) \\ &+ \frac{m_D^2 - m_K^2}{q^2}q^\alpha f_0(q^2) \end{aligned} \tag{2.20}$$

$$< K(k)|\bar{s}\sigma^{\alpha\beta}c|D(p) > = im_D^{-1}f_2(q^2)(p^\alpha k^\beta - p^\beta k^\alpha) \tag{2.21}$$

$$< K(k)|\bar{s}c|D(p) > = \frac{m_D^2 - m_K^2}{m_c - m_s}f_0(q^2) \tag{2.22}$$

where q is defined as $q = p - k$. The amplitude of this decay process is given by

$$\begin{aligned} < \nu(p_1)\bar{l}(p_2)K(k)|\mathcal{H}_{NP}|D_s(p) >= \frac{G_F}{\sqrt{2}}\left(\bar{u}(p_1)(1+\gamma_5)G_V\left(\not{Q}f_+(q^2)\right.\right. \\ + \left(\frac{m_D^2 - m_K^2}{q^2}\right)\left(f_0(q^2) - f_+(q^2)\right)\not{q}\Big)v(p_2) + G_S f_0(q^2)\left(\frac{m_D^2 - m_K^2}{m_c - m_s}\right)\bar{u}(p_1)(1+\gamma_5)v(p_2) \\ \left. - \frac{G_T f_2(q^2)}{2m_D}\bar{u}(p_1)(1+\gamma_5)(\not{q}\not{Q} - \not{Q}\not{q})v(p_2)\right) \end{aligned} \tag{2.23}$$

where $Q = p+k$. The G-functions in the former equation are defined as $G_V = V_{cs}^* + (C_{sc\ell\nu}^{V,LL} + C_{sc\ell\nu}^{V,RL})/2\sqrt{2}$, $G_S = (C_{sc\ell\nu}^{S,RR} + C_{sc\ell\nu}^{S,LR})/2\sqrt{2}$ and $G_T = (C_{sc\ell\nu}^{T,RR} + C_{sc\ell\nu}^{T,LR})/2\sqrt{2}$. Hence, the amplitude square of the process, writen in terms of the vector, scalar and tensor contributions, is

$$
\begin{aligned}
|A_{D\to k\nu l}|^2 = 4G_F^2 \Bigg[|G_V|^2 B_V + \left|(G_S) f_0(q^2)\left(\frac{m_D^2 - m_K^2}{m_c - m_s}\right)\right|^2 B_S + |G_T f_2(q^2)|^2 B_T \\
+ 2\Re(G_V G_S^*) f_0(q^2)\left(\frac{m_D^2 - m_K^2}{m_c - m_s}\right) B_{VS} + 2\Re(G_V G_T^*) f_2(q^2) B_{VT} \\
+ 2\Re(G_S G_T^*) f_0(q^2)\left(\frac{m_D^2 - m_K^2}{m_c - m_s}\right) f_2(q^2) B_{ST}\Bigg]
\end{aligned}
\tag{2.24}
$$

where the B-functions are:

$$
\begin{aligned}
B_V &= -\Bigg[p_1\cdot p_2\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+)\left((m_D^2 - m_K^2)(f_0 - f_+) + 2f_+ q\cdot Q\right) + |f_+|^2 Q^2\right) \\
&- 2\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+) p_2\cdot q + f_+ p_2\cdot Q\right)\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+) q\cdot p_1 + f_+ Q\cdot p_1\right)\Bigg] \\
B_S &= p_1\cdot p_2 \\
B_T &= \frac{2}{m_D^2}\Big(p_2\cdot q(Q^2 q\cdot p_1 - q\cdot Q Q\cdot p_1) - p_2\cdot Q(q\cdot Q q\cdot p_1 - q^2 Q\cdot p_1) \\
&+ \frac{1}{2} p_1\cdot p_2\left[(q\cdot Q)^2 - q^2 Q^2\right]\Big) \\
B_{VS} &= -m_l\left(\frac{(m_D^2 - m_K^2)}{q^2}(f_0 - f_+) q\cdot\nu + f_+ Q\cdot\nu\right) \\
B_{ST} &= -\frac{1}{m_D}(q\cdot p_1 Q\cdot p_2 - q\cdot p_2 Q\cdot p_1) \\
B_{VT} &= \frac{m_l}{m_D}\Big(q\cdot\nu\left(\frac{(m_D^2 - m_K^2}{q^2}(f_0 - f_+) q\cdot Q + f_+ Q^2\right) \\
&- Q\cdot\nu\left((f_0 - f_+)(m_D^2 - m_K^2) + f_+ q\cdot Q\right)\Big)
\end{aligned}
\tag{2.25}
$$

Finally, in the rest frame (RF) of the decaying meson, the partial decay rate for the $D^0 \to K^{\pm} l^{\mp} \nu$ decay channel with non standard interactions is given by

$$\begin{aligned} \frac{d\Gamma_{D\to K\ell\nu_\ell}}{dE_K} &= \frac{G_F^2 m_D \sqrt{E_K^2 - m_K^2}}{(2\pi)^3} \Big\{ (E_K^2 - m_K^2) \frac{2q^2 + m_\ell^2}{3q^2} \left| G_V f_+(q^2) \right|^2 \\ &+ \left(-|G_T f_2(q^2)|^2 \frac{q^2 + 2m_l^2}{3} + m_l G_V f_+(q^2) G_T^* f_2(q^2) \right) \left(\frac{E_k^2 - m_K^2}{m_D^2} \right) \\ &+ \frac{|(m_D^2 - m_K^2) q f_0(q^2)|^2}{4m_D^2} \left| \frac{m_\ell}{q^2} G_V + \frac{G_S}{m_c - m_s} \right|^2 \Big\} \left(1 - \frac{m_\ell^2}{q^2} \right)^2 , \end{aligned} \quad (2.26)$$

Other constants involved in Eqs. (2.19, 2.25) are: G_F the Fermi constant, V_{cs}^* the CKM matrix element, $m_\ell, m_c, m_s, m_K, m_{D_s}, m_D$ the masses of the leptons, charm and strange quarks, the Kaon and D meson respectively as reported by PDG [7]. The transferred energy is $q^2 = m_D^2 + m_K^2 - 2m_D E_K$ and E_K is the final energy of the Kaon meson. Its allowed energy is $m_K < E_K < (m_D^2 + m_K^2 - m_\ell^2)/2m_D$. The decay constant f_{D_s} in the leptonic decay rate is defined by $\langle 0|\bar{s}\gamma_\mu\gamma_5 c|D_s(p)\rangle = i f_{D_s} p_\mu$. In the semileptonic decays, the scalar, vector and tensor form factors $f_0(q^2)$, $f_+(q^2)$ and $f_2(q^2)$ are defined via $\langle K|\bar{s}\gamma^\mu c|D\rangle = f_+(q^2)(p_D + p_K - \Delta)^\mu + f_0(q^2)\Delta^\mu$, with $\Delta^\mu = (m_D^2 - m_K^2)q^\mu/q^2$, $\langle K|\bar{s}c|D\rangle = (m_D^2 - m_K^2)/(m_c - m_s) f_0(q^2)$ and $\langle K(k)|\bar{s}\sigma^{\mu\nu} c|D(p)\rangle = i m_D^{-1} f_2(q^2)(p^\mu k^\nu - p^\nu k^\mu)$.

2.3.3 Transverse polarization including NSIs

The transverse polarization of the charged lepton in the decay $D \to Kl\nu$ is a sensitive T-violating or CP violating observable when CPT is conserved. This observable was first computed in the semileptonic decay $K^+ \to \pi^0 \mu^+ \nu$ as a useful tool for studying non standard CP violation [58, 59] . In the SM, P_T is expected to be highly suppressed, as in the case of the charged Kaon $K_{\mu 3}$ [60] or neutral Kaon K^0 [61]. Given the similarities with the K^+ decay, we can compute the transverse polarization for the semileptonic decay of the D meson

$D(p)^0 \to K^{\mp}(k)\nu(p_1)l(p_2)^{\pm}$. The transverse polarization is given by

$$P_T^S = \frac{|A_T^S|^2 - |A_T^{-S}|^2}{|A_T^S|^2 + |A_T^{-S}|^2} \tag{2.27}$$

where S represents the spin of the lepton. In general, one measures the spin perpendicular to the decay plane defined by the final particles [50]. Thus in order to have a non zero effect the transverse polarization should be proportional to $\epsilon^{\alpha\beta\gamma\delta}p_\alpha p_{1\beta}p_{2\gamma}S_\delta$, where p_1 and p_2 are the 4-vectors of neutrino and charged lepton respectively. Given that $S_\mu = (0, \mathbf{s})^T$, with s perpendicular to the decay plane, the polarized amplitude can be written as

$$\begin{aligned}
|A_T^S|^2 &= \frac{1}{2}|A_{D\to K\ell\nu_\ell}|^2 \\
&+8G_F^2\epsilon^{\alpha\beta\gamma\delta}K_\alpha S_\beta p_{1\gamma}p_{2\delta}\Bigg[f_+(q^2)f_0(q^2)\frac{M_D^2 - M_K^2}{m_c - m_s}\mathrm{Im}(G_V G_S^*) \\
&+ f_2(q^2)\Bigg((f_0(q^2) - f_+(q^2))(M_D^2 - M_K^2)(1 - 2\frac{p_2\cdot q}{q^2}) \\
&+ f_+(q^2)\frac{q\cdot Q - 2p_2\cdot Q}{M_D}\Bigg)\mathrm{Im}(G_V G_T^*) \\
&+ f_2(q^2)f_0(q^2)\frac{m_\ell}{M_D}\frac{M_D^2 - M_K^2}{m_c - m_s}\mathrm{Im}(G_T G_S^*)\Bigg],
\end{aligned} \tag{2.28}$$

here $Q = p + k$ and $q = p - k$. With this, we can construct the transverse polarization averaged over the charged lepton energy. To calculate the averaged transverse polarization we have integrated over the charged lepton energy. Thus this observable can be written in the decay frame of the D meson as

$$\begin{aligned}
\langle P_T^S\rangle &= \frac{G_F^2 M_D^4}{4\pi^3}\left(\frac{d\Gamma}{dE_K}\right)^{-1}\Bigg\{f_0(q^2)f_+(q^2)\frac{M_D^2 - M_K^2}{M_D(m_c - m_s)}g_0(q^2)\mathrm{Im}(G_V G_S^*) \\
&+ f_2(q^2)f_0(q^2)\frac{m_\ell}{M_D^2}\frac{M_D^2 - M_K^2}{m_c - m_s}g_0(q^2)\mathrm{Im}(G_T G_S^*) \\
&+ f_2(q^2)\Bigg[f_0(q^2)\frac{m_\ell^2}{M_D^2}\frac{(M_D^2 - M_K^2)}{q^2}g_0(q^2) + f_+(q^2)\Big(g_1(q^2) \\
&- \frac{m_\ell^2 + q^2}{M_D^2}\frac{(M_D^2 - M_K^2 - q^2)}{q^2}g_0(q^2)\Big)\Bigg]\mathrm{Im}(G_T G_V^*)\Bigg\}
\end{aligned} \tag{2.29}$$

where we have defined the dimensionless kinematical functions

$$g_0(q^2) \equiv \frac{1}{M_D^3} \int_{E_\ell^{\rm min}}^{E_\ell^{\rm max}} dE_\ell |\mathbf{p_1} \times \mathbf{p_2}| \,, \qquad g_1(q^2) \equiv \frac{1}{M_D^4} \int_{E_\ell^{\rm min}}^{E_\ell^{\rm max}} dE_\ell E_\ell \cdot |\mathbf{p_1} \times \mathbf{p_2}| \,, \tag{2.30}$$

with

$$E_\ell^{\rm max\ (min)} = \frac{1}{2}(m_D - E_K)\left(\frac{m_\ell^2}{q^2} + 1\right) \pm \frac{1}{2}\sqrt{E_K^2 - m_K^2}\left(1 - \frac{m_\ell^2}{q^2}\right) . \tag{2.31}$$

We can see that the leading contributions in New Physics are the scalar and tensor interactions, i.e. at first order in C's.

2.4 Model independent analysis and experimental constraints

2.4.1 D meson decays measurements *vs* theoretical branching ratios

In this work we make a combined analysis of the D meson leptonic and semileptonic decays (fig. 2.3) to constrain the generic Wilson coefficients that parameterize NP effects in the $c\bar{s} \to l\bar{\nu}$ transition. We use the decays,

- $D^0 \to K^- l^+ \nu$, $l = e, \mu$
- $D^+ \to K^0 l^+ \nu$, $l = e, \mu$
- $D_s^+ \to l^+ \nu$, $l = \mu, \tau$

There are a number of measurable observables related to the D meson that might be modified by NSI. D_s leptonic decays have been measured by a number of experiments, namely CLEO [64] and Belle [65] among other experiments. Semileptonic decays, on the other hand, have been observed with an integrated luminosity of 818pb^{-1} [68, 69, 70]. In particular, the q^2 distribution for the semileptonic decays $D^+ \to \bar{K}^0 e^+ \nu_e$, $D^0 \to K^- e^+ \nu_e$ has been measured

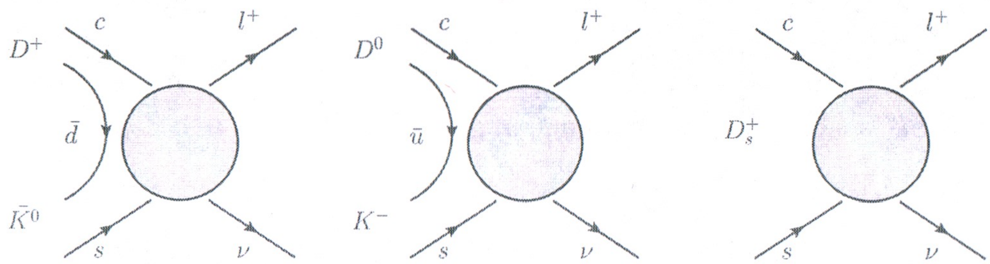

Figure 2.3: $c\bar{s} \to \ell\bar{\nu}$ transition in D meson leptonic and semileptonic decays.

Table 2.1: Theoretical and experimental branching ratios

i	Decay	Theoretical BR $\mathcal{B}_i^{th}$	Experimental BR $\mathcal{B}_i^{exp}$
1	$D^0 \to K^- e^+ \nu_e$	$(3.28 \pm 0.11)\%$.	$(3.55 \pm 0.04)\%$
2	$D^0 \to K^- \mu^+ \nu_\mu$	$(3.22 \pm 0.11)\%$	$(3.30 \pm 0.13)\%$
3	$D^+ \to \bar{K}^0 e^+ \nu_e$	$(8.40 \pm 0.32)\%$.	$(8.83 \pm 0.22)\%$
4	$D^+ \to \bar{K}^0 \mu^+ \nu_\mu$	$(8.24 \pm 0.31)\%$	$(9.2 \pm 0.6)\%$
5	$D_s^+ \to \tau^+ \nu_\tau$	$(5.10 \pm 0.22)\%$	$(5.43 \pm 0.31)\%$
6	$D_s^+ \to \mu^+ \nu_\mu$	$(5.20 \pm 0.20) \times 10^{-3}$	$(5.90 \pm 0.33) \times 10^{-3}$

by CLEO [66],[67]. From those measurements it is possible to extract the lifetimes for the mesons. They result to be $\tau_{D^0} = (410.1 \pm 1.5) \times 10^{-15}$ s, $\tau_{D^+} = (1040 \pm 7) \times 10^{-15}$ s, and $\tau_{D_s} = (500 \pm 7) \times 10^{-15}$ s. In summary, total branching ratios for semileptonic decays of the D^0 and D^+ and the world measured total branching ratios for the leptonic decays of the D_s are shown in Table 2.1.

We compute the theoretical decay rates, on the other hand, $\Gamma^{th}_{D_s \to \ell\nu}$, $\Gamma^{th}_{D^0 \to K^+ \ell^- \nu_\ell}$ and $\Gamma^{th}_{D^+ \to \bar{K}^0 \ell^+ \nu_\ell}$ given by eqs. (2.19,2.25) fixing all the Wilson's coefficients to zero. The resulst are shown in table 2.1. We ignore all radiative corrections since they are expected to be below the 1% [62].

Other relevant physical inputs needed for the SM computation of the theoretical BRs are:

1. *The CKM element V_{cs}.* As we are looking for New Physics, we have to be very careful on the value of the CKM element we will use in our numerical analysis. In order to avoid that leptonic and semileptonic of D mesons have been used to fix the V_{cs} value, we use the central value of the CKM element which comes from $W \to cs$ decay, neutrino-nucleon scattering and unitary constraints coming from $b - s$ transitions relating $|V_{cd}|$ and $|V_{cs}|$ through unitarity. This last constraint gives the strongest constraint. So our central value for V_{cs} is $|V_{cs}| = 0.97344 \pm 0.00016$ [63]. Using this unitary constraint means that automatically our results will not apply to any model with more than three fermion families.

2. *Hadronic form factors.* These are non-perturbative parameters calculated in specific theoretical models. In particular Lattice QCD is a well-established method able to compute the hadronic form factors from first principles, that has reached an excellent precision [6]. Therefore, for our analysis, we fix the hadronic form factors and leptonic decay constant to the value estimated with lattice QCD simulations. The leptonic decay constant f_{D_s}, defined as $\langle 0|\bar{s}\gamma_\mu\gamma_5 c|D_s(p)\rangle = i f_{D_s} p_\mu$, has been computed with a precision of the order of 2% by the HPQCD collaboration [39]. In order to compute the leptonic branching ratio we have used the reported value of $f_{D_s} = 248 \pm 2.5$ MeV [39]. On the other hand, less is known about $f_0(q^2), f_+(q^2)$ which are defined via $\langle K|\bar{s}\gamma^\mu c|D\rangle = f_+(q^2)(p_D + p_K - \Delta)^\mu + f_0(q^2)\Delta^\mu$, with $\Delta^\mu = (m_D^2 - m_K^2)q^\mu/q^2$, and $\langle K|\bar{s}c|D\rangle = (m_D^2 - m_K^2)/(m_c - m_s) f_0(q^2)$. Dramatic progress has been made over the last decade on lattice calculations for those form factors [39, 40, 6]. We use the latest results by the HPQCD collaboration [6] as input for the calculation of the theoretical

decay rate.

We use the same functional shape of the form factors as in [66]:

$$f_+(q^2) = \frac{f_+(0)}{(1-\frac{\alpha q^2}{M_X^2})(1-\frac{q^2}{M_X^2})}$$
$$f_0(q^2) = \frac{f_+(0)}{(1-\frac{q^2}{M_X^2\beta})} \tag{2.32}$$

where M_X is the appropiate pole mass, i.e. $M_{D_s^*}$ for f_+ and $M_{D_{s0}^*}$ for f_0; $\alpha = 0.5$ and $\beta = 1.31$ [66]. At $q^2 = 0$, $f_+(0) = 0.745 \pm 0.011$ [6]. We stress the fact that using theoretical form factors, calculated from first principles, e.g. from Lattice QCD is more effective and accurate when constraining non standard interactions. In order to illustrate this point we show a quantitive analysis varying a single Wilson coefficient together with the parameters $f_+(0), \alpha, \beta$. Using the experimental data of the q^2 dependance for the decay rates $D^+ \to K^0 e^+ \nu_e$, $D^0 \to K^- e^+ \nu_e$ [67], and the experimental branching ratios (table 2.1) with their corresponding uncertainties we show the allowed regions at 90% and 68% c.l. for $C_1 = C^{VLL}_{sce\nu_e}$ and $f_+(0)$ in figure 2.4. Note that $f_+(0)$ is less accurate than the value calculated from Lattice QCD [6], as expected. It is important to stress that the decay rate is proportional to the square of the form factors, therefore their uncertainty is expected to be relevant, in fact it is the main source of uncertainty in our calculations.

The results for the theoretical BRs are listed in table 2.1. with their corresponding uncertainties. The total theoretical uncertainties are calculated straightforward: propagating each uncertainty for every physical constant as reported in PDG[7], and the theoretical uncertainties coming from the lattice QCD calculations of the form factors. The main contribution in the theoretical error comes from the leptonic decay constant f_{D_s} and the semileptonic form factors $f_+(q^2)$ and $f_0(q^2)$. The reported error in f_{D_s} induces a $\sim 4\%$ error in the theoretical

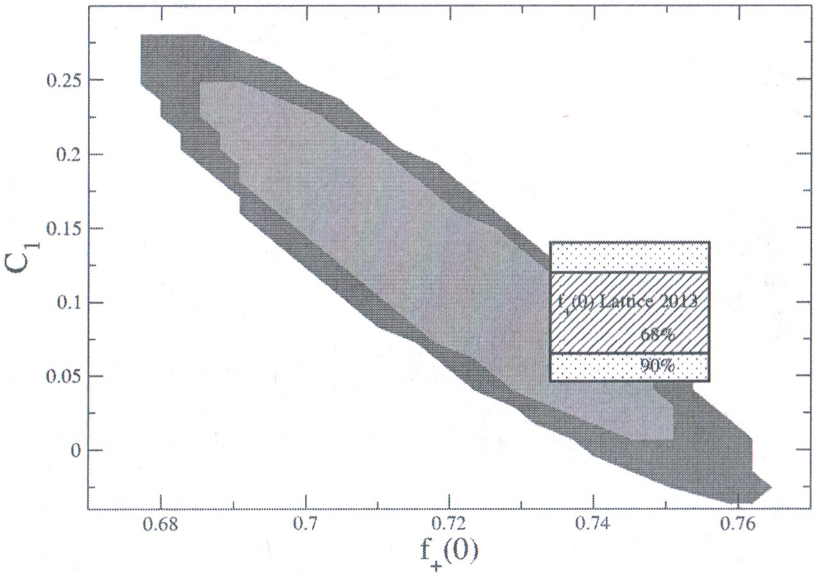

Figure 2.4: Allowed regions at 90% (cyan) and 68% (red) c.l. for $C_1 = C^{VLL}_{sce\nu_e}$ and $f_+(0)$ from experimental data compared to Lattice QCD results [6]

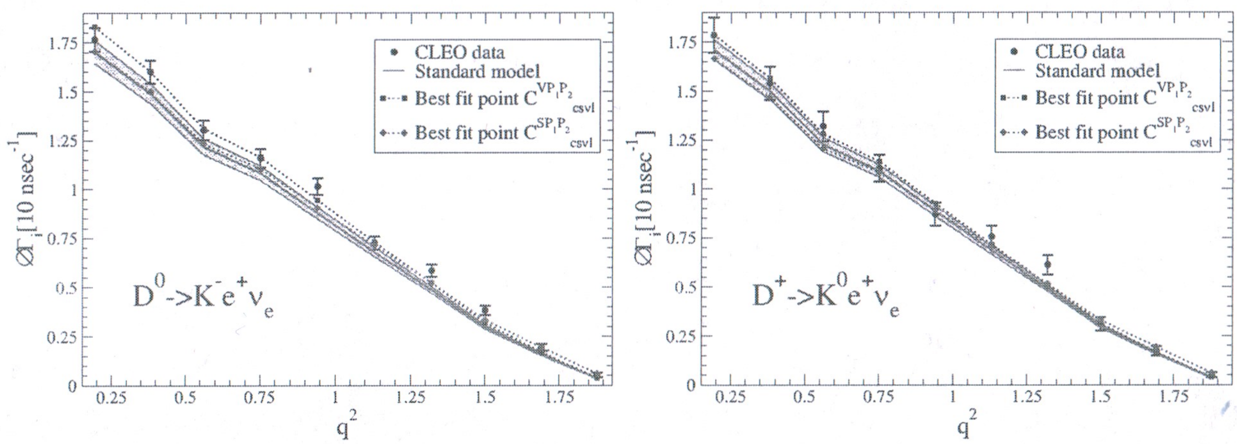

Figure 2.5: Partial decays measured by CLEO [67] and the theoretical partial decay computed with the Standard model using the latest form factors from [6]. Grey region represents one sigma theoretical error. Good agreement is observed

	4 pars. 95% C.L.	χ^2_{min}/d.o.f	1 par. 95% C.L.	χ^2_{min}/d.o.f
$C^{V,LL}_{sc\ell\nu}$	$[-.094, 0.42]$	0.62	$[0.072, 0.14]$	0.89
$C^{V,RL}_{sc\ell\nu}$	$[-0.34, 0.17]$	0.62	$[0.057, 0.13]$	1.29
$C^{S,RR}_{sc\ell\nu}$	$[-0.33, 0.21]$	0.62	$[-0.22, -0.21] \cup [0.00, 0.13]$	2.19
$C^{S,LR}_{sc\ell\nu}$	$[-0.23, 0.33]$	0.62	$[-0.012, 0.00] \cup [0.20, 0.22]$	2.17

Table 2.2: Model independent constraints at 95% C.L. for universal non standard interactions using leptonic and semileptonic D meson decays. We have fixed the leptonic decay constant and semileptonic form factors to those estimated by lattice QCD. In the first column, four parameters are allowed to vary at a time and in the third column, only one parameter is varied.

leptonic branching ratio. Similarly, the reported error in the lattice determination of $f_0(q^2)$ and $f_+(q^2)$ leads to a $\sim$ 4% error in the theoretical semileptonic branching ratio. Exact values are listed in table 2.1. As already mentioned, world average measurements of the total BRs as reported by PDG [7] are shown in Table 2.1 for comparison. In the same way, the theoretical partial decays for the $D^0 \to K^+e^-\nu_e$ and $D^+ \to \bar{K}^0e^+\nu_e$ and the CLEO data points are shown in Figure 2.5. Note the good agreement between experiment and theory.

The BRs reported in table 2.1 are a pure theoretical prediction of the SM in the following sense: $\mathcal{B}^{th}$ is computed using the SM Lagrangian only, since we have set all Wilson coefficients to zero, and the form factors are computed from first principles using Lattice results[6].

2.4.2 Constraining real NSI

Let us assume that the new physics effects, are parameterized, as described in section 2.3, by the Wilson coefficients. In this first part of our analysis we suppose the non standard physical phases are aligned with those of the SM in such a way that in general we can consider the Wilson coefficients real. We compute the range of the Wilson coefficients to exactly match the theory and the experiment. In order to do so, we perform a simple χ^2 analysis, with

Flavor dependent scalar non standard interactions

	95% C.L. $\Re^+$	$\frac{\chi^2_{min}}{\text{d.o.f}}$	95% C.L. $\Re$	$\frac{\chi^2_{min}}{\text{d.o.f}}$
$C^{S,RR}_{sce\nu_e} + C^{S,LR}_{sce\nu_e}$	$[0.32, 0.47]$	1.05	$[-0.47, -0.33] \cup [0.32, 0.47]$	1.05
$C^{S,RR}_{sc\mu\nu_\mu}$	$[0.0, 0.27]$	1.17	$[-0.77, 0.25]$	1.30
$C^{S,LR}_{sc\mu\nu_\mu}$	$[0.0, 0.38]$	1.17	$[-0.63, 0.38]$	1.30
$C^{S,RR}_{sc\tau\nu_\tau} - C^{S,LR}_{sc\tau\nu_\tau}$			$[-0.075, 0.175]$	1.0

Flavor dependent vector non standard interactions

	95% C.L. $\Re^+$	$\chi^2_{min}/\text{d.o.f}$	95% C.L. $\Re$	$\chi^2_{min}/\text{d.o.f}$
$C^{V,LL}_{sce\nu_e} + C^{V,RL}_{sce\nu_e}$	$[0.07, 0.14]$	0.99	$[0.07, 0.14]$	0.99
$C^{V,LL}_{sc\mu\nu_\mu}$	$[0.0, 0.22]$	1.67	$[-0.025, 0.255]$	0.93
$C^{V,RL}_{sc\mu\nu_\mu}$	$[0.0, 0.1]$	1.67	$[-0.19, 0.095]$	0.93
$C^{V,LL}_{sc\tau\nu_\tau} - C^{V,RL}_{sc\tau\nu_\tau}$			$[-0.12, 0.28]$	1.0

Table 2.3: Model independent constraints at 95% C.L. for scalar and vector flavor dependent non standard interactions from the leptonic and semileptonic D meson decays. We have fixed the leptonic decay constant and semileptonic form factors to those estimated by lattice QCD. In the second column, the Wilson coefficients are restricted to be positive. In the fourth column, the Wilson coefficients are only restricted to be real numbers.

$\chi^2 = \sum_i (\mathcal{B}_i^{th} - \mathcal{B}_i^{exp})^2/\delta\mathcal{B}_i^2$. Here, $\delta\mathcal{B}_i$ is calculated adding in quadratures the experimental and theoretical uncertainties shown in Table 2.1.

We shall consider first a combined analysis of the leptonic and semileptonic BRs and the experimental data from CLEO assuming only scalar (S) and vector (V) NSI. An analysis including all the New Physics operators at a time, scalar, vector and tensor, shows that the tensor contribution is negligible as compared to the former operators. However we can constrain the tensor interactions assuming that only the tensor operator is dominant, as we show in the next subsection. Hence, the relevant parameters with the above considerations are: $C_{sc\ell\nu}^{V,LL}, C_{sc\ell\nu}^{V,RL}, C_{sc\ell\nu}^{S,RR}$ and $C_{sc\ell\nu}^{S,LR}$. Although this is a restrictive hypothesis, this analysis is useful for models where no CP violating phases or models in which the phyiscal phases are aligned with the CKM phase, e.g THDM-II or some specific MSSM-$\not{R}$ as we will show later. The results for the relevant Wilson coefficients, assuming these are flavor universal or flavor dependent, are shown in tables 2.2 and 2.3 respectively.

- *Flavor independent NSI* Table 2.2 corresponds to universal NSI, that is, flavor independent interactions. When we do not take into account the tensor interaction, we are left with four coefficients: $C_{sc\ell\nu}^{V,LL}, C_{sc\ell\nu}^{V,RL}, C_{sc\ell\nu}^{S,RR}$ and $C_{sc\ell\nu}^{S,LR}$. Notice that equations (2.19,2.25) have a different dependence on the Wilson coefficients, hence, when combining the leptonic decay rates and the semileptonic decay rates it is possible to extract a bound for each parameter even if we analyze the four parameters at a time. We have computed the allowed values for those universal coefficients at 95% C.L. by varying the four parameters at-a-time, i.e. those are the most general cases, this is because both scalar and vector universal NSI may affect the Brs. On the other hand, we have also estimated the allowed regions by varying only one parameter at a time (right column Table 2.2), this is when only one lepton flavor independent NSI contributes to the physical process.

- *Flavor dependent NSI* Some models may induce only vector, as well as only scalar NSI at a time. As we will show in the next section, the left-right model or the two Higgs doublet model are examples of each type of NSI, respectively. In those cases, we can obtain the bounds for the corresponding Wilson coefficients. Those coefficients may depend on the flavor of the lepton involved. Since we have only six $\mathcal{B}_i^{th}$s, we can perform the χ^2 analysis only if we assume scalar NSI or vector NSI at a time. In each case, for the electron NSI, we use the channels $i = 1, 3$ and the CLEO data points from the kinematic distribution, for the muon $i = 2, 4, 6$ and for the tau, only a fit can be performed with $i = 5$; channel i as shown in Table 2.1. Results for both cases, scalar and vector flavor dependent NSI are listed table 2.3. As we have mentioned, those constraints can be applied to the THDMs. In those cases, Wilson coefficients are positive. Hence, we have constrained Wilson coefficients either assuming they are real positive numbers or just real numbers. We will show the effectiveness of those constraints for specific models.

2.4.3 Complex Wilson coefficients

We shall consider now complex flavor universal Wilson coefficients. Many models of New Physics introduce CP violating phases which are in general not aligned with the SM CP violating phase, therefore we also analyze such scenario. Here, we assume that only one non-standard operator is dominant besides the Standard Model operator, either scalar, vector or tensor NSIs. This means we will take into account only one complex Wilson coefficient at a time, i.e. two independent parameters for each operator. We consider again a combined analysis of the leptonic and semileptonic BRs and the experimental data from CLEO.

The model independent constraints at 68%, 90% and 95% confidence level are shown in

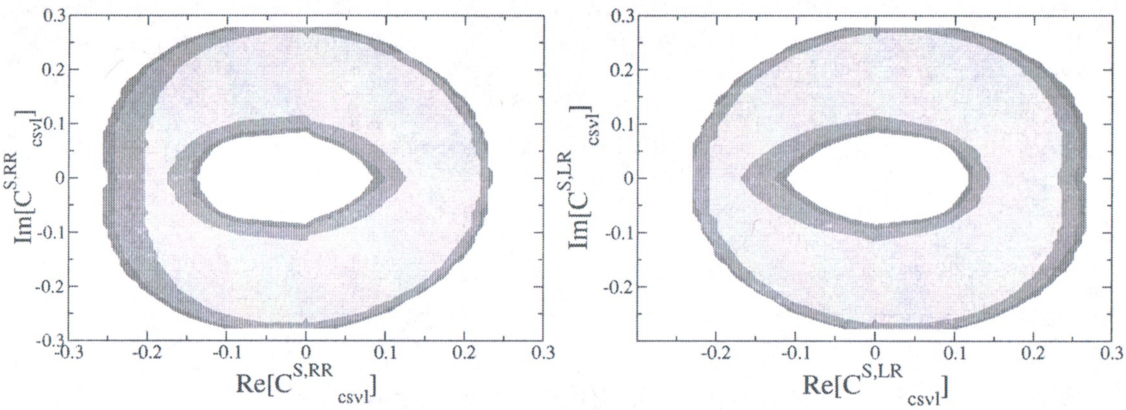

Figure 2.6: Universal scalar NSI, parametrized by the complex coefficients $C^{S,rr}_{sc\ell\nu}, C^{S,LR}_{sc\ell\nu}$, allowed from D meson decays. Colored regions correspond to 68% , 90% and 95% C.L respectively

Figures 2.8, 2.6, 2.7. Contrary to the scalar or vector NSI, tensor NSI can not be separated from the unknown form factor $f_T(0) \equiv f_2(0)$. Hence, we can only obtain the bounds for $Re[f_T G_T]$ and $Im[f_T G_T]$, shown in Figure 2.7. In summary, the allowed regions at 95% C.L. are the following:

- vector NSI: $\chi^2/\text{d.o.f.} = 0.96$

$$\begin{aligned}
-0.5 < \mathrm{Re}[C^{V,LL}_{sc\ell\nu}] < 0.21, &\qquad 95\%\ \text{C.L.}\,, \\
-1.63 < \mathrm{Im}[C^{V,LL}_{sc\ell\nu}] < 1.63, &\qquad 95\%\ \text{C.L.}\,, \\
-0.9 < \mathrm{Re}[C^{V,RL}_{sc\ell\nu}] < 0.7, &\qquad 95\%\ \text{C.L.}\,, \\
-2.1 < \mathrm{Im}[C^{V,RL}_{sc\ell\nu}] < 2.1, &\qquad 95\%\ \text{C.L.}\,.
\end{aligned} \tag{2.33}$$

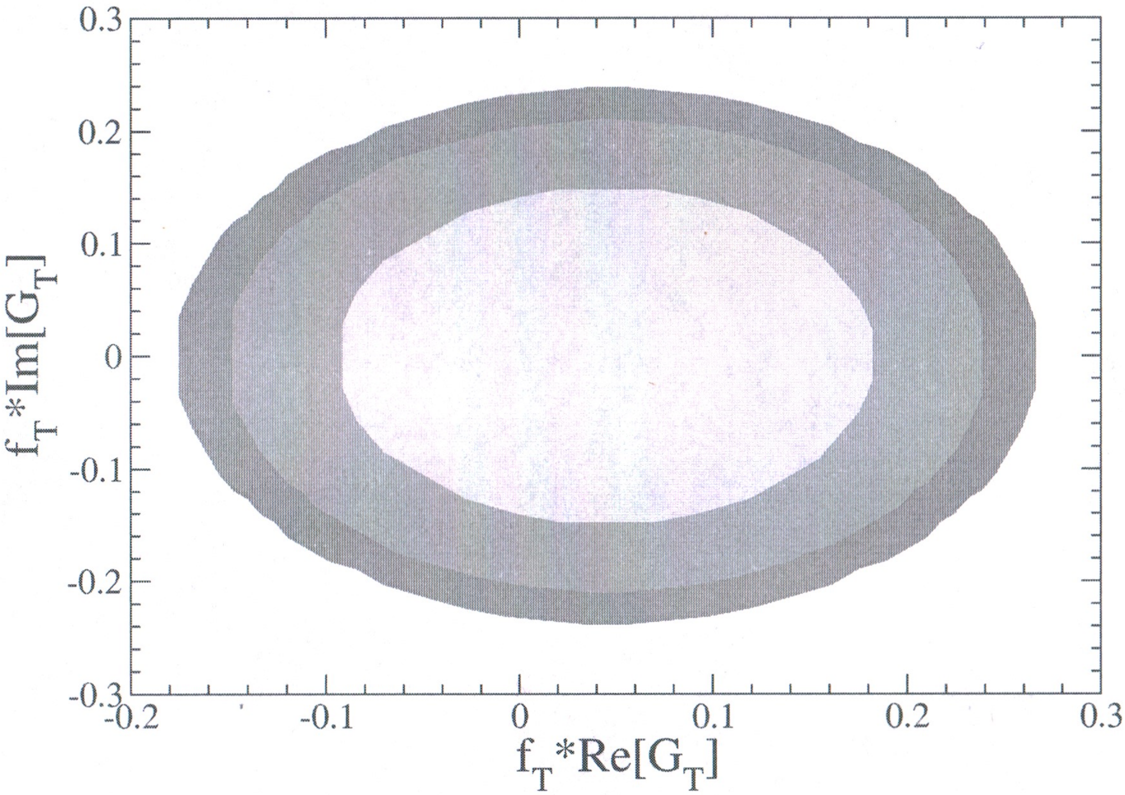

Figure 2.7: Contrary to scalar or vector NSI, tensor NSI can not be separated from the unknown form factor $f_T(0)$. Hence, we can only obtain the respective bounds for $Re[f_T G_T]$ and $Im[f_T G_T]$. Colored regions correspond to 68% , 90% and 95% C.L respectively

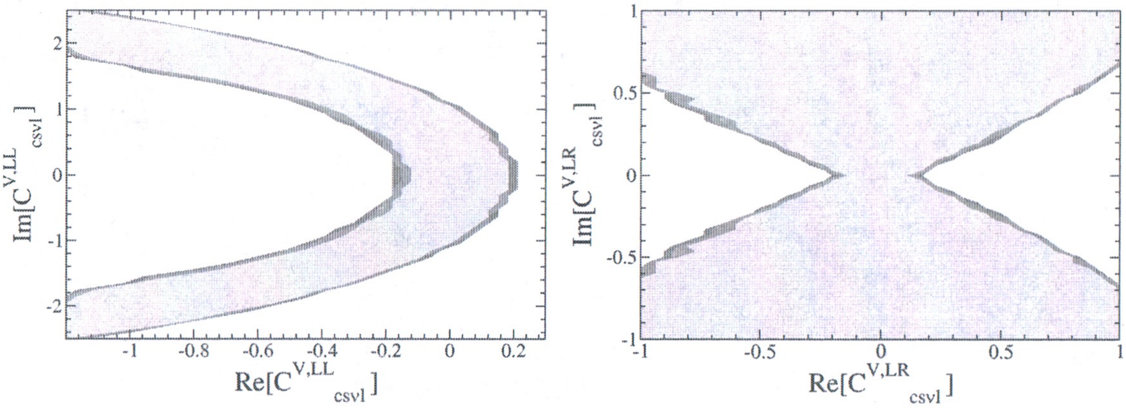

Figure 2.8: Universal vector NSI, parametrized by the complex coefficients $C^{V,LL}_{sc\ell\nu}$, $C^{V,RL}_{sc\ell\nu}$, allowed from D meson decays. Colored regions correspond to 68% , 90% and 95% C.L respectively

- Scalar NSI: $\chi^2/\text{d.o.f.} = 1.20$

$$
\begin{aligned}
-0.24 < \mathrm{Re}[C^{S,RR}_{sc\ell\nu}] < 0.23, &\qquad 95\%\ \text{C.L.}, \\
-0.28 < \mathrm{Im}[C^{S,RR}_{sc\ell\nu}] < 0.28, &\qquad 95\%\ \text{C.L.}, \\
-0.23 < \mathrm{Re}[C^{S,LR}_{sc\ell\nu}] < 0.26, &\qquad 95\%\ \text{C.L.}, \\
-0.29 < \mathrm{Im}[C^{S,LR}_{sc\ell\nu}] < 0.29, &\qquad 95\%\ \text{C.L.}. \qquad (2.34)
\end{aligned}
$$

We use the best fit points to compute the partial decays of the D meson, $D^+ \to \bar{K}^0 e^+ \nu_e$ and $D^0 \to K^- e^+ \nu_e$, and we show them in Figure 2.5, compared with the experimental data and the Standard Model prediction. For those points we see there is better agreement with the experimental data.

2.4.4 Transverse polarization estimation

As an application of our results we give a prediction for a T-odd observable, the transverse polarization of the charged lepton for the decay $D^+ \to \bar{K}^0 \ell^+ \nu_\ell$. This observable has not been measured. We chose this semileptonic decay thinking the experimental measurement could be done as in the case of the K^+ meson, [50]. The K^+ decays as $K^+ \to \pi^0 \ell^+ \nu_\ell$ and the BR of the $\pi^0 \to \gamma\gamma$ is $BR(\pi^0 \to \gamma\gamma) = 98.823 \pm 0.034\%$ [7], this allows for a clean distinction of the angular distribution of the charged lepton, hence the transverse polarization. In our case, the $\bar{K}^0$ decays with a BR of $BR(\bar{K}^0 \to \pi^0\pi^0) = 30.69 \pm 0.05\%$ [7], allowing possibly for a distinction of the angular distribution of the charged lepton. In the SM, P_T is expected to be highly suppressed, as in the case of the charged Kaon $K_{\mu 3}$ [60] or neutral Kaon K^0 [61]. This implies that a large value, i.e. $P_T \gtrsim \mathcal{O}(10^{-3})$ is a signal of new physics. As we performed the analysis for the complex universal Wilson coefficients taking into account only one dominant non-standard operator the transverse polarization (2.28) can only be computed for each case. Furthermore, notice that if in the the transverse polarization (2.28) we only take into account the vector contribution it will vanish. For these reason we show the only non-vanishing transverse polarizations including New Physics integrated over all the kinematical allowed region. The results are shown in Figure (2.9).

We can see in Figure (2.9) that there is little dependence on the real contributions of both scalar and tensor non-standard interactions. The largest value of the transverse polarization allowed from the previous constraints over the complex universal Wilson coefficients is $P_T = 0.23$, which is not negligible.

2.5 Model Dependent Analysis:

Let us consider now different models of New Physics. We perform a χ^2 analysis in a model dependent way by finding the respective bounds over the relevant parameters for those models. In particular, we obtain bounds for the Two Higgs Doublet Model Type-II and Type III, the Left-Right model, the Minimal Supersymmetric Standard Model with explicit R-Parity violation and Leptoquarks. We show that under some simplifying assumptions, the model

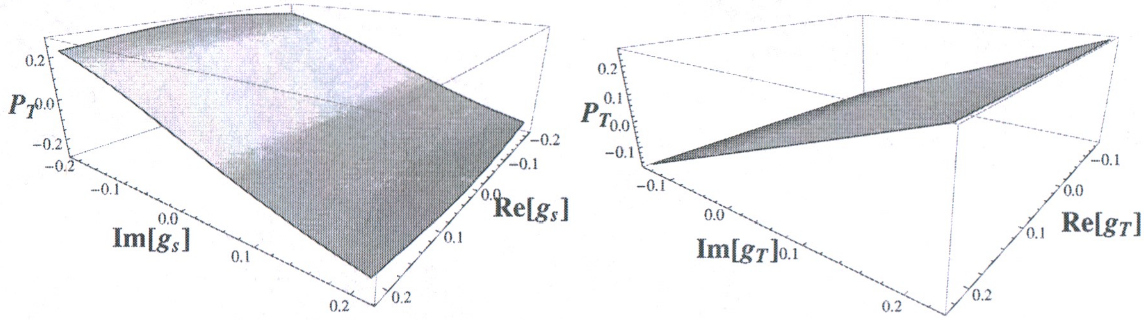

Figure 2.9: Estimated transverse polarization P_T for the $D^+ \to \bar{K}^0 \ell^+ \nu_\ell$. the left figure shows the P_T when only scalar non-standard interactions are considered. The right figure is the P_T when only tensor non-standard interactions are considered. The largest value of P_T allowed from the previous constraints over the complex universal Wilson coefficients is $P_T = 0.23$

independent constraints can be mapped to some particular models, exemplifying the usefulness of this kind of analysis.

2.5.1 Two Higgs doublet model (THDM):

It is one of the simplest and economical extensions of the SM, see[71, 72] for a review. THDM introduces an additional scalar doublet that induces scalar charged currents ($H^\pm$), two neutral scalar fields and a pseudoscalar neutral field (h^0, H^0 and A^0). For D meson decays, the only two parameters involved are the new scalar mass (m_{H^+}) and the ratio of the vacuum expectation values $\tan\beta$ of the two Higgs doublets. At low energies, the Lagrangian for THDM, in the Higgs basis for the charge scalars and the mass basis for fermions, is given by [73]

$$\begin{aligned} -\mathcal{L}_{H^\pm} &= \sqrt{2}/v H^+ [V_{u_i d_j} \bar{u}_i (m_{u_i} X P_L + m_{d_j} Y P_R) d_j + \\ &+ m_\ell Z \bar{\nu}_L \ell_R] + \text{H.c.} \end{aligned} \tag{2.35}$$

with X, Y, Z functions of m_{H^+} and $\tan\beta$ different for different versions of THDM, and the Wilson coefficients will be given by

$$\frac{C^{S,RR}_{sc\ell\nu}}{2\sqrt{2}} = V^*_{cs}\frac{m_\ell m_c}{M_H^2}ZX\,, \quad \frac{C^{S,LR}_{sc\ell\nu}}{2\sqrt{2}} = V^*_{cs}\frac{m_\ell m_s}{M_H^2}ZY\,. \tag{2.36}$$

In particular, THDM-II has Natural Flavour Conservation, namely the suppression of Flavor Changing Neutral Currents (FCNC) at tree level, through a Z_2 symmetry [74]. Interesting bounds have been obtained with meson decay experiments [75] and recently LEP has reported a lower bound on the mass of the charged Higgs of 80 GeV [76].

For THDM-II, $X = \cot\beta$, $Y = Z = \tan\beta$. We perform a χ^2 analysis using the 26 observables from the leptonic and semileptonic BRs and the kinematical distribution from CLEO. The result is shown in Figure 2.10. We can see from this figure that D meson decays favor lower masses for the charged Higgs at 90% C.L., however at 95%, there is good agreement with LEP bounds.

Now we will illustrate the effectiveness of our model independent bounds, once we apply them to Wilson coefficients of THDM, eqs (2.35). There is a flavor dependence coming from the mass of the leptons involved. Since this is an scalar interaction, we can use the bounds on flavor dependent scalar NSI. From $C^{S,RR}_{sc\tau\nu_\tau} - C^{S,LR}_{sc\tau\nu_\tau}$ we get the region -1.8×10^{-3} GeV^{-1} $<$ $(m_c - m_s\tan^2\beta)/M_H^2 < 0.023$ GeV^{-1} at 68% C.L., which gives the outer region of an ellipse and the inner region of an hyperbole in the plane $(m_H, \tan\beta)$ illustrated in Fig. 2.10. Those regions are in excellent agreement with the region obtained by a complete χ^2 analysis performed with all D meson decays. The allowed values for $\tan\beta$ and m_H^+ are plotted in Fig. 2.10 in a shadow gray area. This agreement illustrates the effectiveness of using generic Wilson coefficient to constrain the relevant parameters of models beyond the SM. A more complete analysis for the THDM-II model using different observables from flavor physics has been done in [77]. Our aim in this work is not to compete with those constraints, rather than illustrate the effectiveness of this type of generic analysis and to show that semileptonic

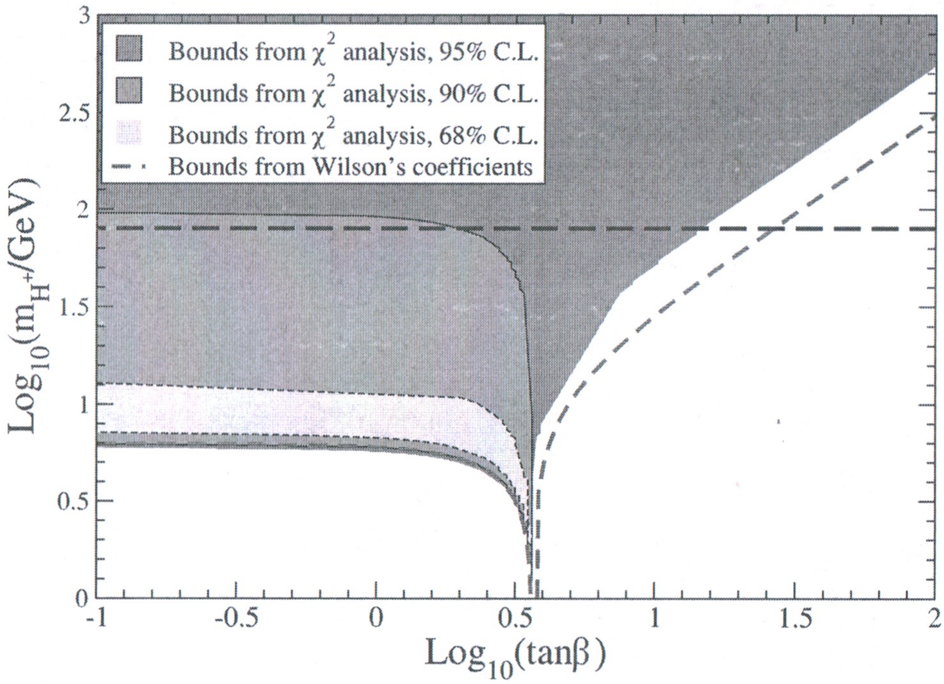

Figure 2.10: Allowed regions for $\tan\beta$ and the mass of the charged Higgs to be consistent with the D meson decays at 68% C.L., 90% C.L. and 95% C.L. obtained performing a complete χ^2 analysis of the BRs. Dashed lines are the limits at 90% C.L. using the bounds on Wilson coefficients (Table 2.3) showing good agreement. As a reference, the LEP limit on the mass of a charged Higgs is also plotted [76].

decays may shed complementary information.

For completeness, let us briefly mention the THDM-III case which can be analyzed immediately by noting that the Wilson coefficients in this case correspond to the following definitions:

$$X = \cot\beta - \frac{\csc\beta}{\sqrt{2\sqrt{2}G_F}} m_c^{-1} \left(\tilde{Y}^u_{1,22} + \frac{V_{us}}{V_{cs}} \tilde{Y}^u_{1,21} + \frac{V_{ts}}{V_{cs}} \tilde{Y}^u_{1,23} \right), \tag{2.37}$$

$$Y = \tan\beta - \frac{\sec\beta}{\sqrt{2\sqrt{2}G_F}} m_s^{-1} \left(\tilde{Y}^d_{2,22} + \frac{V_{cd}}{V_{cs}} \tilde{Y}^d_{2,12} + \frac{V_{cb}}{V_{cs}} \tilde{Y}^d_{2,32} \right), \tag{2.38}$$

$$Z = \tan\beta - \frac{\sec\beta}{\sqrt{2\sqrt{2}G_F}} m_\ell^{-1} \tilde{Y}^\ell_{2,\nu_\ell \ell} \tag{2.39}$$

$$\tag{2.40}$$

where $\tilde{Y}^f_{a,ij}$ are the Yukawa elements as were defined in [78, 79]. The corresponding bounds obtained via eqs. 2.35 for THDM-III are interesting since they show relations between $\tilde{Y}^f_{a,ij}$, β and the mass of the charged Higgs.

2.5.2 Left-right model:

As an example of a model with vector NSI as the main contribution to new physics, we will consider SM's extensions based on extending the SM gauge group including a gauge $SU(2)_R$. The original model, based on the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, restores the parity symmetry at high energies [42, 43, 44, 45, 46]. This SM extension has been extensively studied in previous works (see for instance refs.[80, 81, 82, 83, 84]). Recent bounds on the mass of W_R[7, 85, 86, 87, 88, 89, 90, 91, 92, 93] have strongly constrained these models. TWIST collaboration [94, 95] found a model independent limit on ξ to be smaller than 0.03 (taking $g_L = g_R$ at muon scale) through precision measurements of the muon decay parameters. However the presence of right-handed currents may weaken some tensions between inclusive and exclusive determinations of some CKM elements [96, 97, 98], so it is appropriate to explore less restrictive versions of the LR model. Recently and ample phenomenological analysis has been done for the LR model using B physics [99], nevertheless we shall see that current D physics can shed complementary bounds on the free parameters of the model for a specific scenario. Here we consider the scenario where Left-Right is not manifest, that is $g_L \neq g_R$ at unification scale, with the presence of mixing between left and right bosons through a mixing angle ξ. This LR mixing is restricted by deviation to non-unitarity of the CKM quark mixing matrix. In case of manifest LR model, it is well known that ξ has to be smaller than 0.005[100] and $M_{W'}$ bigger than 2.5 TeV[101]. But in the no manifest case, the constraint on $M_{W'}$ are much less restrictive as $M_{W'}$ could be as light as

0.3 TeV[102] and ξ can be as large as 0.02[82, 103, 104, 105]. The Lagrangian for this case, including only the vertex of interest, is given by

$$\begin{aligned}
-\mathcal{L}_{LR} &= \frac{g_L}{\sqrt{2}}\bar{u}_i\gamma^\mu\left[\left(c_\xi V_{u_id_j}P_L + \frac{g_R}{g_L}s_\xi \bar{V}^R_{u_id_j}P_R\right)W^+_\mu\right. \\
&\left. + \left(-s_\xi V_{u_id_j}P_L + \frac{g_R}{g_L}c_\xi \bar{V}^R_{u_id_j}P_R\right)W'^+_\mu\right]d_j \\
&+\frac{g_L}{\sqrt{2}}\bar{\nu}_L\gamma^\mu\left(c_\xi W^+_\mu - s_\xi W'^+_\mu\right)\ell_L + \text{H.c.}, \qquad (2.41)
\end{aligned}$$

where $c_\xi = \cos\xi$ and $s_\xi = \sin\xi$ and W^+, W'^+ are the mass states of gauge bosons. Likewise $\bar{V}^R_{u_id_j} = \exp^{-i\omega} V^R_{u_id_j}$, where ω is a CP violating phase. After integrating degrees of freedom in a usual way, this leads to the Wilson coefficients

$$C^{V,LL}_{sc\ell\nu} = \sin^2\xi\left(\frac{M^2_W}{M^2_{W'}} - 1\right)V_{cs} \qquad (2.42)$$

$$C^{V,RL}_{sc\ell\nu} = \frac{g_L}{2g_R}\sin(2\xi)\left(1 - \frac{M^2_W}{M_{W'}}\right)\bar{V}^R_{cs} \qquad (2.43)$$

Such scenarios were studied for instance in [106]. In our case, the relevant parameters are: ξ, $M_{W'}, g_L/g_R\text{Re}[V^R_{cs}]$ and $g_L/g_R\text{Im}[V^R_{cs}]$. By performing the combined analysis for all our 26 observables, by varying these four parameters at a time we found the allowed regions for ξ and $M_{W'}$ which are shown in Fig. 2.11. There is only one viable restriction for the following parameters: $-71.0 < g_L/g_R\text{Re}[V^R_{cs}] < 83$, while the analysis is insensitive to the imaginary part.

2.5.3 MSSM-$\not{R}$:

R-Parity is a discrete symmetry defined as $(-1)^{3B+L+2S}$, where B, L and S are the baryon number, lepton number and particle spin respectively. R-Parity violating (RPV) interactions involve either lepton number violation or baryon number violation, but not both in order to preserve proton stability. These interactions lead to flavor violating interactions in the

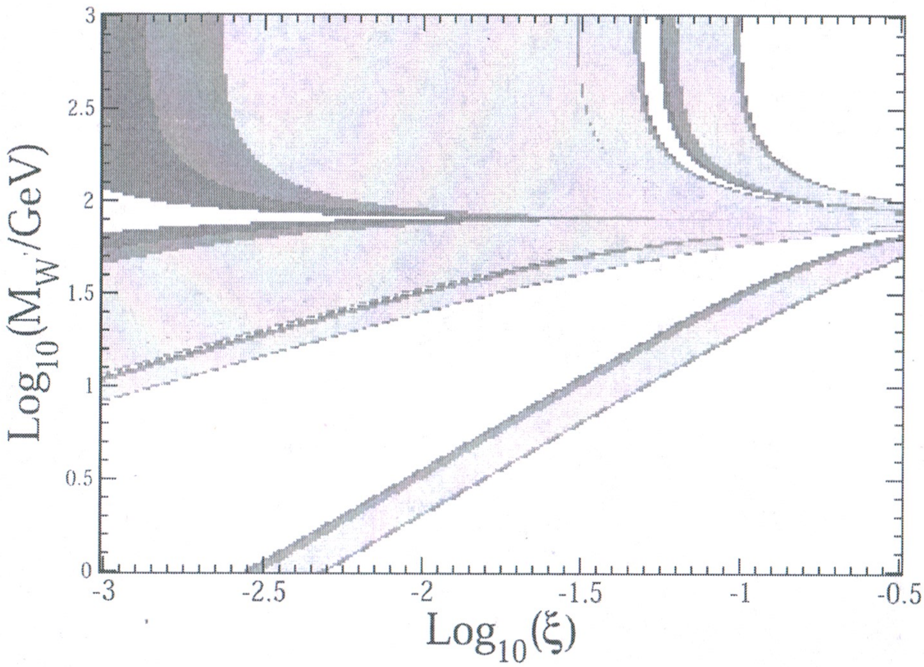

Figure 2.11: Bounds at 68% ,90% and 95% C.L. on ξ and the mass of the W' boson obtained by using D meson decays data.

leptonic and hadronic sector, and read explicitly as,

$$\begin{aligned}
\mathcal{L}_S &= \lambda_{ijk}\left[\tilde{\nu^i}_L\overline{e^k_R}e^j_L + \tilde{e^j}_L\overline{e^k_R}\nu^i_L + \tilde{e^{*k}}_R\overline{(\nu^i_L)^c}e^j_L - (i \leftrightarrow j)\right] \\
&+ \lambda'_{ijk}\left[\tilde{\nu}^i_L\overline{d^k_R}d^j_L + \tilde{d}^j_L\overline{d^k_R}\nu^i_L + \tilde{d^*}^k_R\overline{(\nu^i_L)^c}d^j_L - \tilde{e}^i_L\overline{d^k_R}u^j_L\right. \\
&- \left.\tilde{u}^j_L\overline{d^k_R}e^i_L - \tilde{d^*}^k_R\overline{(e^i_L)^c}u^i_L\right] + h.c. \qquad (2.44)
\end{aligned}$$

A vast majority of observables have been used to set the corresponding bounds to these effective couplings (for a complete review see [29] and references therein); in particular, for D meson decays [107, 108, 109, 110, 111, 112, 113, 114, 115, 116]. We relate the corresponding Wilson coefficients constrained in the global analysis to the RPV couplings which constructively interfere with the Standard Model, i.e. through the exchange of a $-1/3$ electrically charged squark in a t-channel, which fixes the neutrino flavor, described by,

$$\begin{aligned}
\mathcal{L}_S &= V^*_{cs}\frac{\sum_k |\lambda'_{i2k}|^2}{m^2_{\tilde{d}^{k*}_R}}(\bar{\nu_L}^i s^c_R \bar{l^c_R} c_L) \\
&\xrightarrow{\text{Fierz}} V^*_{cs}\frac{\sum_k |\lambda'_{i2k}|^2}{2m^2_{\tilde{d}^{k*}_R}}(\bar{\nu_L}^i\gamma^\mu l^i_L \bar{s_L}\gamma_\mu c_L), \qquad (2.45)
\end{aligned}$$

where a Fierz transformation is done to rearrange the former operator in terms of the product of a leptonic and a hadronic current. The only non-vanishing Wilson coefficient is

$$C^{V,LL}_{sc\ell\nu} = \sqrt{2}V_{cs}/G_F\sum_k |\lambda'_{i2k}|^2/m^2_{\tilde{d}^{k*}_R} \qquad (2.46)$$

Using the conservative bounds for the model independent constraints (table 2.3) we get the following constraints at 95% confidence level and expressed in GeV^{-2}:

$$\begin{aligned}
0.05 < \sum_k |\lambda'_{12k}|^2/(m_{\tilde{d}^{k*}_R}/300\text{GeV})^2 &< 0.11, \\
\sum_k |\lambda'_{22k}|^2/(m_{\tilde{d}^{k*}_R}/300\text{GeV})^2 &< 0.17, \\
\sum_k |\lambda'_{32k}|^2/(m_{\tilde{d}^{k*}_R}/300\text{GeV})^2 &< 0.22 \qquad (2.47)
\end{aligned}$$

Our bounds agree with those found in [113] for muon and tau flavor. Nevertheless it is interesting to note that for the electron flavor we find more restrictive bounds. This is taking into account the latest and more accurate values of the form factors from lattice QCD as previously mentioned.

2.5.4 The effective leptoquark lagrangian

Leptoquark particles are scalars or vectors bosons that carry both baryon number and lepton number [47],[48]. These new particle states are expected to exist in various extensions of the SM. Leptoquark states usually emerge in grand unified theories (GUTs) [117, 118, 119] (as vector) or technicolor models (as scalars) [120, 121], or SUSY models with R-parity violation (as we saw in the previous section), but are described naturally in low energy theories as an effective for fermion interaction of a more fundamental theory. Effective interactions induced by leptoquark exchange can be manifest in meson decays, in particular, for the second generation of quarks in D meson decays. A vast majority of observables have been used to set the corresponding bounds to these effective couplings; in particular, for D meson decays [122, 52, 123]. Leptoquarks are usually classified by their appropiate quantum numbers under the gauge symmetries of the Standard Model, such as colour, hypercharge, and isospin charge [48]. These particles may couple to both quark chiralities, the left handed or right handed, in particular the scalar leptoquark S. When we rearrange the effective interactions in order to have external quark and lepton currents we do some Fierz transformations that lead to tensor, scalar and vector interactions, that we shall take into account in a model dependent analysis. We will consider the exchange of the scalar leptoquarks: S_0 with charge $-1/3$ and $(3, 1, -2/3)$ gauge numbers; and the $S_{1/2}$ with charge $2/3$ and $(3, 2, 7/3)$ gauge numbers.

Hence, the effective Lagrangian for the $c \rightarrow s$ transition (Fig.2.1) is given by

$$\begin{aligned} L^{LQ}_{Eff} &= V^*_{cs} \left[\frac{\kappa^{R*}_{i2} \kappa^{L}_{i2}}{m^2_{S^{2/3}_{1/2}}} (\overline{\nu^i_L} c_R \overline{l^c_{iL}} s^c_R) \right. \\ &- \left. \frac{\kappa'^{R*}_{i2} \kappa'^{L}_{i2}}{m^2_{S^{-1/3}_0}} (\overline{\nu^i_L} s^c_R \overline{l^c_{iL}} c_R) - \frac{|\kappa'^{L}_{i2}|^2}{m^2_{S^{-1/3}_0}} (\overline{\nu^i_L} s^c_R \overline{l^c_{iR}} c_L) \right] \end{aligned} \tag{2.48}$$

After a Fierz transformation we have, in terms of the Wilson operators,

$$\begin{aligned} \mathcal{L}^{LQ}_{Eff} = \frac{1}{2} V^*_{cs} & \left[\left(\left(\frac{\kappa^{R*}_{i2} \kappa^{L}_{i2}}{m^2_{S^{2/3}_{1/2}}} + \frac{\kappa'^{R*}_{i2} \kappa'^{L}_{i2}}{m^2_{S^{-1/3}_0}} \right) (\overline{\nu_L}^i l_{iR} \overline{s_L} c_R \right. \right. \\ & - \left. \left. \frac{1}{4} \overline{\nu_L}^i \sigma_{\mu\nu} l_{iR} \overline{s_L} \sigma^{\mu\nu} c_R \right) + \frac{|\kappa'^{L}_{i2}|}{m^2_{S^{-1/3}_0}} (\overline{\nu}^i \gamma^\mu P_L l_i \overline{s} \gamma_\mu P_L c) \right] \end{aligned} \tag{2.49}$$

Note that $C^{VLL}_{sc\ell\nu}, C^{VLR}_{sc\ell\nu}, C^{SRR}_{sc\ell\nu} = -4 C^{TRR}_{sc\ell\nu}$ given by,

$$C^{TRR}_{sc\ell\nu} = \frac{\sqrt{2} V_{cs}}{G_F} \left(\frac{\kappa^{R*}_{i2} \kappa^{L}_{i2}}{m^2_{S^{2/3}_{1/2}}} + \frac{\kappa'^{R*}_{i2} \kappa'^{L}_{i2}}{m^2_{S^{-1/3}_0}} \right), \quad C^{VLL}_{sc\ell\nu} = \frac{\sqrt{2} V_{cs}}{G_F} \left(\frac{|\kappa'^{L}_{i2}|^2}{m^2_{S^{-1/3}_0}} \right) \tag{2.50}$$

Where the last Wilson coefficient also derives from the SUSY $\not{R}$ Lagrangian. In the following we show the respective bounds as a result from our χ analysis considering the 26 observables: the leptonic and semileptonic decays of the D meson and the CLEO data points of the q^2 distribution. Notice here that we have one complex and one real independent Wilson coefficients (as the tensor operator is proportional to the scalar operator), and the tensor form factor f_T. However this analysis is not sensitive to the tensor form factor as the tensor contribution is negligible when the scalar and vector interactions are taken into account, which are the dominant contributions. Hence the model dependent analysis is done varying 3 parameters at a time. At 95% C.L. and expressed in GeV^{-2} these are given by:

$$-0.17 < \mathrm{Re}\left(\kappa_{i2}^{R*}\kappa_{i2}^{L} + \kappa_{i2}^{\prime R*}\kappa_{i2}^{\prime L}\right)/(m_S/300\mathrm{GeV})^2 < 0.01,$$
$$-0.09 < \mathrm{Im}\left(\kappa_{i2}^{R*}\kappa_{i2}^{L} + \kappa_{i2}^{\prime R*}\kappa_{i2}^{\prime L}\right)/(m_S/300\mathrm{GeV})^2 < 0.10,$$
$$0.04 < |\kappa_{i2}^{\prime L}|^2/(m_{S_0}/300\mathrm{GeV})^2 < 0.11 \qquad (2.51)$$

As an example we can consider the leptoquark states that couple to the second generation of left handed quarks (chiral generation leptoquarks) and the first generation of left handed leptons. This means we take into account only the coefficient in Eq. (2.49), which also derives from the SUSY $\not{R}$ effective lagrangian and corresponds to Eq. (2.45). Therefore the Wilson coefficient is real and flavor dependent on the first generation of leptons, hence, we use the model independent constraints obtained in Section (2.4) for flavor dependent parameters, given in Table (2.3) which correspons to the first constraint in Eq. (2.46). The allowed region at 95% C.L. from the semileptonic decays of the D mesons is given by:

$$0.05\left(\frac{m_{S_0}}{300\mathrm{GeV}}\right)^2 < |\kappa_{12}^{\prime}|^2 < 0.11\left(\frac{m_{S_0}}{300\mathrm{GeV}}\right)^2 \qquad (2.52)$$

Previous bounds [7] for the second generation of left handed quarks couplking to the first generation of left handed leptons, are reported to be $\kappa'^2 < 5 \times (M_{LQ}/300\mathrm{GeV})^2$ for S_0. As stated in the previous subsection (for the MSSM-$\not{R}$), for the electron flavor and the second generation of quarks, this former constraint is more restrictive than previous bounds.

Chapter 3

Conclusions

Low energy experiments are reaching an outstanding precision and a vast amount of measurable observables. This allow us to do a test for deviations from the Standard Model or test fundamental properties with a high level of accuracy. In this work we showed two interesting methods to search (indirectly) for New Physics in the low energy range, of the order of $\mathcal{O}(\mathrm{GeV})$ within the leptonic and the quark sector. We found encouraging results.

In the first chapter we addressed a fundamental property of the neutrino, its nature, i.e. if it is a Majorana or Dirac fermion. If Majorana neutrinos do exist, $|\Delta L| = 2$ processes like neutrino-antineutrino oscillations can occur. The production of leptons with same charges at the production and detection vertices of neutrinos will be a clear manifestation of these processes. In this work we have used the S-matrix formalism of quantum field theory to describe these oscillations in the case of muon neutrinos produced in π^+ decays which convert into muon antineutrinos that are detected via inverse beta decay on nucleons.

One interesting result is that the time evolution probability of the whole process is not factorizable into production, oscillation and detection probabilities, as is the case in neutrino oscillations [27]. We find that, for very short times of propagation of neutrinos, the observation of $\mu^+\mu^+$ events would lead to a direct bound on the effective mass of muon Majorana neutrinos. In the case of long-baseline neutrino experiments, the CP rate asymmetry for production of $\mu^+\mu^+/\mu^-\mu^-$ events would lead to direct bounds on the difference of CP-violating Majorana phases. Finally, using the current bound on muon neutrino-antineutrino oscillations reported by the MINOS Collaboration we are able to set the bound $\langle m_{\mu\mu}\rangle \lesssim 64$ MeV, which is the first direct limit on the neutrino masses, although it is still several orders of magnitude below current indirect bounds reported in the literature.

Future results from MINOS are expected from the analysis of twice the data set used to get the bound reported so far [26] and quoted in Eq. (12) above. Since current uncertainties in the observed and expected number of $\bar{\nu}_\mu$ events are dominated by statistical errors [26], we could expect only a slight improvement by a factor $1/\sqrt{2}$ on the effective Majorana mass of the muon neutrino. Neutrino factories may improve this bound by more than one order of magnitude. As a consequence of these results, neutrino experiments aiming to measure neutrino-antineutrino oscillations with different short- and long-baseline setups can be useful to get direct and complementary constraints on the masses and phases of Majorana neutrinos.

On the other hand, we assume that the hypothetic observation of $\Delta L = 2$ processes in neutrino-antinutrino oscillations is due to non standard interactions, at first order, in the production of the neutrino. As an example we analysed the R-parity violating superpotential of the MSSM, for the case in which the observation of the final states, same-charged muons at the production and detection of neutrinos, is a result of non standard lepton number violating interactions at the neutrino source.

In the second chapter of this thesis we analyse the case for non standard interactions in the leptonic and semileptonic decays of the D meson. These processes have been measured with high precision and have not yet fully exploited. We showed how can we constrain generic parametrizations of non-standard interactions only with D meson decays, which resulted to be sufficiently restrictive.

We have combined the $D_s \to \ell\nu_\ell$ and the semileptonic $D^0 \to K\ell\nu_\ell$ and $D^+ \to K\ell\nu_\ell$ decays, together with the q^2 distribution of the semileptonic decays for the electron channel measured by CLEO. The theoretical BRs were computed with the latest lattice results on $f_+(q^2)$ and $f_0(q^2)$ form factors [6]. We have found the corresponding bounds for the Wilson coefficients that parameterize the contribution of new physics as non standard interactions. We considered two scenarios in which the New Physics models have either aligned or not aligned physical phases with those of the SM, i.e. real or complex Wilson coefficients. Those constraints can be applied to some model of new physics generating scalar, vector, or tensor operators, such as the THDM-II, Type III, the Left-Right model, MSSM-$\not{R}$ and leptoquarks, which we analyzed here. We show the usefulness of the model independent constraints as well as specific cases when a model dependent analysis is needed. In our model dependent analysis we found that for the THDM-II a low mass for the charged Higgs is favored, at 90% C.L. 6.3GeV $< m_{H^+} <$ 63.1GeV. We showed there are no strong restrictions for the LR model with these combined decay channels but comparable with previous bounds [102, 82, 103, 104, 105]. In particular for the MSSM-$\not{R}$, our bounds agree with those found in [113] for muon and tau flavor. Nevertheless it is interesting to note that for the electron flavor we find more restrictive bounds. For the leptoquark model, taking into account the couplings to the second generation of left handed quarks and first generation of left handed leptons the constraints coming from D meson decays are $0.05\left(\frac{m_{S_0}}{100\text{GeV}}\right)^2 < |\kappa'_{12}|^2 < 0.11\left(\frac{m_{S_0}}{100\text{GeV}}\right)^2$. These results are encouraging, at least ten times less, if compared with previous bounds $|\kappa'_{12}|^2 < 1.7 \times (m_{s_0}/100\text{GeV})^2$ [7]. This is taking into account the latest and more accurate values of the form factors from lattice QCD as previously mentioned

We estimated the transverse polarization and found that these model independent constraints obtained from the D meson decays allow a large P_T, which is expected to be highly suppressed in the SM. Hence the experimental measure of P_T could be useful to constraint Wilson coefficients involving quarks of the second generation.

Appendix A

Appendix

A.1 $\Delta L = 2$ processes in neutrino oscillations

A.1.1 Kinematic allowed region for $\pi p \to n\mu\mu$

The amplitud square in terms of invariant quantities is

$$\begin{aligned}
\left|T_{\nu_\mu - \bar{\nu}_\mu}(t)\right|^2 &= (G_F V_{ud} f_\pi)^2 64 m_\nu^2 \left(((g_a^2 - 1)^2 p_1 \cdot p_3 (2p_a \cdot p_2 p_a \cdot p_2 - m_\pi^2 p_b \cdot p_2)\right. \\
&+ ((g_a + 1)((g_a + 1) p_b \cdot p_3 (2p_a \cdot p_2 p_a \cdot p_1 - m_\pi^2 p_1 \cdot p_2) \\
&+ \left.(g_a - 1) m_n m_p (2p_a \cdot p_2 p_a \cdot p_3 - m_\pi^2 p_2 \cdot p_3))))\right) \qquad \text{(A.1)}
\end{aligned}$$

where $p_a, p_b, p_1, p-2, p_3$ are the four momentum of the pion, proton, neutron, muon 1 (production) and muon 2 (detection) respectively for the process $\pi(p_a)p_(p_b) \to n(p_1)\mu(p_2)\mu(p_3)$, which is a $\Delta L = 2$ process. The total cross section is given by

$$\begin{aligned}
\sigma(s) &= \left(\frac{\pi}{(32(2\pi)^5 \lambda(s, m_\pi^2, m_p^2))} \frac{1}{\sqrt{-\Delta}} (G_F V_{ud} f_\pi)^2 64 m_\nu^2 \left(((g_a^2 - 1)^2 p_1 \cdot p_3 (2p_a \cdot p_2 p_a \cdot p_2\right.\right. \\
&- m_\pi^2 p_b \cdot p_2) + ((g_a + 1)((g_a + 1) p_b \cdot p_3 (2p_a \cdot p_2 p_a \cdot p_1 - m_\pi^2 p_1 \cdot p_2) \\
&+ \left.\left.(g_a - 1) m_n m_p (2p_a \cdot p_2 p_a \cdot p_3 - m_\pi^2 p_2 \cdot p_3))))\right)\right) \qquad \text{(A.2)}
\end{aligned}$$

where we define $\lambda(x, y, z) = (x - y - z)^2 - 4yz$ and $\Delta = det|\Delta_{ij}|$, and Δ_{ij} is [124]

$$\Delta_{ij} = \begin{pmatrix} 2m_\pi^2 & s-m_\pi^2-m_p^2 & m_\pi^2+m_n^2-t_1 & s-s_1+t_2-m_p^2 \\ s-m_\pi^2-m_p^2 & 2m_p^2 & s-s_2+t_1-m_\pi^2 & m_p^2+m_l^2-t_2 \\ m_\pi^2+m_n^2-t_1 & s-s_2+t_1-m_\pi^2 & 2m_n^2 & s-s_1-s_2+m_l^2 \\ s-s_1+t_2-m_p^2 & m_p^2+m_l^2-t_2 & s-s_1-s_2+m_l^2 & 2m_l^2; \end{pmatrix} \tag{A.3}$$

The cross section for $2 \to 3$ particle states is described by 5 independent kinematic variables, e.g. the s, s_1, s_2, t_1, t_2 invariants defined as:

$$\begin{aligned} s = (p_a+p_b)^2,\ \ s_1 = (p_1+p_2)^2,\ \ s_2 = (p_2+p_3)^2 \\ t_1 = (p_a-p_1)^2\ \ t_2 = (p_b-p_3)^2 \end{aligned} \tag{A.4}$$

Finally we can write the lorentz invariants $p_i \cdot p_j$, $i, j = a, b, 1, 2, 3$ in terms of the former invariant quantities,

$$\begin{aligned} p_a \cdot p_b = (s - m_\pi^2 - m_p^2)/2,\ \ p_a \cdot p_1 = (m_\pi^2 + m_n^2 - t_1)/2, \\ p_a \cdot p_2 = (s_1 + t_1 + t_2 - m_n^2)/2,\ \ p_b \cdot p_1 = (s - s_2 + t_1 - m_\pi^2)/2, \\ p_2 \cdot p_3 = (s_2 - m_\mu^2 - m_l^2)/2,\ \ p_a \cdot p_3 = (s - s_1 + t_2 - m_p^2)/2 \\ p_b \cdot p_3 = (m_p^2 + m_l^2 - t_2)/2,\ \ p_1 \cdot p_2 = (s_1 - m_n^2 - m_\mu^2)/2, \\ p_1 \cdot p_3 = (s - s_1 - s_2 + m_\mu^2)/2,\ \ p_b \cdot p_2 = (s_2 + t_2 - t_1 - m_l^2)/2 \end{aligned} \tag{A.5}$$

The kinematic allowed region can not be solved analitically in closed form, but we use a random event generator, ROOT [125], to compute the total cross section in the laboratory system, where the proton is at rest. For low neutrino energies and hence low pion energies it results aproximetly $\sigma \sim 6 \times 10^{-28}\mathrm{GeV}^{-2}$, which is similar to our previous results, making reasonable assumptions over the kinematic range.

A.1.2 CP assymetry

Let us assume that the amplitude of a physical process has two different contributions of the form

$$A = A_1 e^{i\delta_1} + A_2 e^{i\delta_2} \tag{A.6}$$

where $A_{1,2}$ are complex amplitudes and $\delta_{1,2}$ are two CP conserving dynamical phases. The amplitude for the CP conjugate is

$$\bar{A} = A_1^* e^{i\delta_1} + A_2^* e^{i\delta_2} \neq A^* \tag{A.7}$$

In practice we measure the CP asymmetry defined as:

$$a \equiv \frac{\bar{\Gamma} - \Gamma}{\bar{\Gamma} + \Gamma} = \frac{|A|^2 - |\bar{A}|^2}{|A|^2 + |\bar{A}|^2} \tag{A.8}$$

therefore,

$$a = \frac{-2\Im(A_1 A_2^*)\sin(\delta_1 - \delta_2)}{|A_1|^2 + |A_2|^2 + 2\Re(A_1 A_2^*)\cos(\delta_1 - \delta_2)} \tag{A.9}$$

We can conclude that the CP asymmetry is different from cero if the total amplitude has at least two different contributions and there is a relative phase that explicitly violates CP in the theory and a dynamical phase that conserves CP [126].

In neutrino oscillations the dynamical phase comes from the neutrino propagation $e^{-i\frac{m^2 L}{2E}}$, the CP violating phases come from the SM through the PMNS matrix. The CP asymmetry in standard neutrino oscillations is

$$a_{SM} \simeq \frac{-\sum_{i>j} \Im(U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^*)\sin(\triangle m_{ij}^2 \frac{L}{2E})}{2\sum_{i>j} \Re(U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^*)\sin^2(\triangle m_{ij}^2 \frac{L}{4E})} \tag{A.10}$$

A.2 New Physics in D meson decays

A.2.1 D meson semileptonic decay

Here, a different approach, to simplify the kinematics, is followed [128]. The leptonic and hadronic currents are

$$\begin{aligned} L^\mu &= \bar{u}(p_1)\gamma^\mu(1-\gamma^5)v(p_2) \\ H^\mu &=< K(k)|\bar{s}\gamma^\mu c|D(p)> \end{aligned} \tag{A.11}$$

For a pseudoscalar meson in the final state the hadronic current can be written in terms of standard form factors as $< K|\bar{s}\gamma^\mu c|D >= f_+(q^2)(p+k)^\mu + f_-(q^2)(p-k)^\mu$. After summing over the electron and neutrino spins one finds that the amplitude square is of the form $|A(D(p) \to K(k)\ell(p_2)\nu(p_1))|^2 = \frac{1}{2}|G_F V_{cs}|^2 L^{\mu\nu} H_\mu H_\nu^\dagger$ where

$$\begin{aligned} L^{\mu\nu} &= \bar{u}(p_1)\gamma^\mu(1-\gamma^5)v(p_2)\bar{v}(p_2)(1+\gamma^5)\gamma^\nu u(p_1) \\ &= 4(-2g^{\mu\nu}p_1 \cdot p_2 + 2p_1 \cdot p_2 + 2p_2^\mu p_1^\nu) - ie^{\mu\nu\alpha\beta}p_{2\alpha}p_{1\beta} \end{aligned} \tag{A.12}$$

In the center of mass frame of the $\ell\nu$, $\vec{q} = (\vec{p}-\vec{k}) = \vec{p_2}+\vec{p_1} = 0$, and assuming the produced meson is much heavier than the produced leptons (i.e. massless electron) $L^{\mu\nu}$ is given by

$$L^{ij} = 16E_l^2((\delta^{ij} - e^i e^j) + i1/2e^{ijk}e^k) \tag{A.13}$$

where e^i are the unitary vectors of the $\ell\nu$ plane.

It is convinient to define the quantities $y = q^2/m_D^2$ and $x = p \cdot k/m_D^2$, where $q^2 = (p-k)^2$. Neglecting the electron mass the kinematical limits for y are $0 < y < (1-m_K/m_D)^2$. In terms of these parameters $E_l = E_\nu = m_D\sqrt{y}2$. Furthermore, the norm for the vector of the product meson, in the $\ell\nu$ center of mass frame, is given by $|\vec{k}|^2 = (\frac{m_D}{2\sqrt{y}}(1 - m_K^2/m_D^2 - y))^2 - m_K^2 \equiv (\kappa/\sqrt{y})^2$.

We only need then the spatial components of H_μ in the $\ell\nu$ frame. Expanding in terms of the helicity basis we will have

$$\vec{H} = -2\kappa/\sqrt{y} f_+(q^2)\hat{z}, \tag{A.14}$$

and $f_-(q^2)$ makes no contribution as $(p-k)_\mu$ has no spatial component in the $\ell\nu$ frame. Therefore the amplitude square is given by

$$|A(D(p) \to K(k)\ell(p_2)\nu(p_1))|^2 = (G_F)^2 |V_{cs} f_+(q^2)|^2 8 m_D^2 \kappa K^2 \sin^2\theta_e \tag{A.15}$$

The phase space is conviniently written in terms of the $\ell\nu$ center of mass frame (CM) and the frame of the decaying parent meson (DF). In terms of the parameters defined it results $d\Phi = \frac{\kappa m_D}{(4\pi)^5} dy d\Omega_e^{CM} \Omega_K^{DF}$. After integrating over all angles we find that the differential decay rate is

$$\frac{d\Gamma}{dy} = \frac{\kappa^3 m_D^2}{24\pi^3} |G_F V_{cs} f_+(q^2)|^2. \tag{A.16}$$

This former expression is usually found in the literature and is used for the experimental data analysis.

A.2.2 Fierz identities

A Lorentz scalar may be constructed with the 16 bilinear forms in 5 different ways:

$$(\bar{\phi}_1(x)\Gamma\phi_2(x))(\bar{\phi}_3(x)\Gamma\phi_4(x)) \tag{A.17}$$

for $\Gamma = 1_{4\times4}, \gamma^5, \gamma^\mu\gamma^5, \sigma_{\mu\nu}$. Any variant can be expressed as a linear superposition of all the variants with a change of sequence in the spinors.

$$(\bar{\phi}_1\Gamma_i\phi_2)(\bar{\phi}_3\Gamma^i\phi_4) = \sum_k c_{ik}(\bar{\phi}_1\Gamma_k\phi_4)(\bar{\phi}_3\Gamma^k\phi_2) \tag{A.18}$$

where the c_{ik} coefficients are [127]:

	S	V	T	A	P
S	1/4	1/4	-1/4	-1/4	1/4
V	1	1/2	0	-1/2	-1
T	-3/2	0	-1/2	0	-3/2
A	-1	-1/2	0	-1/2	1
P	1/4	-1/4	-1/4	1/4	1/4

In the following list we show the set of Fierz identities used in this work and derived form the former euqtions. The spinor u is the positive energy solution for the Dirac equation (particle) and v is the negative solution (antiparticle).

$$1.\ \bar{u}_1 P_R v_2 \bar{v}_3 P_L u_4 = \frac{1}{2} \bar{u}_1 \gamma^\mu P_L u_4 \bar{u}_2 \gamma_\mu P_L u_3 \tag{A.19}$$

$$2.\ \bar{u}_1 \gamma^\mu P_L u_2 \bar{u}_3 \gamma_\mu P_L u_4 = -\bar{u}_1 \gamma^\mu P_L u_4 \bar{u}_3 \gamma_\mu P_L u_2 \tag{A.20}$$

$$3.\ \bar{u}_1 P_R u_2 \bar{u}_3 P_R u_4 = \frac{1}{2} \bar{u}_1 P_R u_4 \bar{u}_3 P_R u_2 - \frac{1}{8} \bar{u}_1 \sigma_{\mu\nu} P_R u_4 \bar{u}_3 \sigma^{\mu\nu} P_R u_2 \tag{A.21}$$

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Prof. Shaaban Khalil
2014
Director of Center for Theoretical Physics
Zewail City of Science and Technology
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January 14th,

Report on the Ph.D. thesis
"Selected Low Energy Probes For New Physics"
Submitted by
Vannia Gonzalez Macıas

The thesis represents one of the most important and interesting topics in the field of the particle physics phenomenology. It is well presented and very interesting to read. The thesis results were published in 2 journals. This shows the ability of the author to formulate originally a research problem in the field of High Energy Physics, in a highly creative way to analyze this problem in theoretical and numerical context. Therefore, I recommend awarding **Ms. Vannia Gonzalez Macias** the scientific-academic degree **Doctor of Philosophy (PhD)**.

Prof. Shaaban Khalil

León, Gto., a 13 de Enero 2014

Dr. Guillermo Mendoza
Director de la División de Ciencias e Ingenieras

Universidad de Guanajuato

P r e s e n t e

Por este medio, confirmo haber revisado la tesis con titulo " Selected low energy probes for new physics" que la Maestra en Física Vannia Gonzalez Macias esta presentando para obtener el grado de doctor en Física. Despues de la revisión, considero que la tesis cumple con los criterios de calidad pertinentes, de un trabajo de tesis para obtener el grado de Doctor en Física.

Atentamente,

Carlos A. Ramirez

Dr. Carlos A. Ramirez
Departamento de Física
Universidad de los Andes, Bogotá Colombia

Universidad de Guanajuato

León, Gto., a 13 de Enero 2014

Dr. Guillermo Mendoza
Director de la División de Ciencias e Ingenieras

Universidad de Guanajuato

P r e s e n t e

Por este medio, confirmo haber revisado la tesis con titulo " Selected low energy probes for new physics" que la Maestra en Física Vannia Gonzalez Macias esta presentando para obtener el grado de doctor en Física. Despues de la revisión, considero que la tesis cumple con los criterios de calidad que tiene que cumplir un trabajo de tesis para obtener el grado de Doctor en Física.

Atentamente,

Dr. David Yves Ghislain DELEPINE
Profesor Titular B
Departamento de Física

Facultad de Ciencias Físico-Matemáticas

Benemérita Universidad Autónoma de Puebla

Enero, 6, 2014

A quien corresponda.

Por medio de la presente hago de su conocimiento que después de haber leído la tesis de doctorado de la estudiante Vannia González Macías, titulada:

“Selected low-energy probes for new physics”,

considero que la misma cumple con los criterios de originalidad, extensión y correcta presentación de resultados, de modo que se podrá proceder con la programación del examen de grado correspondiente,

Atentamente

Dr. Lorenzo Díaz Cruz
(jldiaz@fcfm.buap.mx)
Profesor Titular, FCFM-BUAP.

Universidad de Guanajuato

León Guanajuato, a; 8 de enero de 2014

Dr. Guillermo Mendoza Díaz
Director de la División de Ciencias e Ingenierías
Campus León, Universidad de Guanajuato
PRESENTE

Estimado Dr. Mendoza:

Me permito informarle que he leído y revisado la tesis titulada **"Selected low energy probes for new physics".** Dicha tesis la realizó la **Maestra en Física Vannia Gonzáles Macias,** como requisito para obtener el grado de Doctor en Física.

Los resultados presentados en la tesis son de interés para la comunidad científica, dado que estudían una amplía variedad de extensiones al modelo estándar de las partículas elementales en distintos procesos; desde neutrinos hasta decaimientos de mesones. Por lo mismo considero que el trabajo hecho por Vannia es de la calidad suficiente para que sea defendida en un exámen profesional, razón por la cual extiendo mi aval para que así se proceda.

Sin más que agregar, agradezco su atención y aprovecho la ocasión para enviarle un cordial saludo.

ATENTAMENTE
"LA VERDAD OS HARÁ LIBRES"

Dr. Juan Barranco Monarca
Departamento de Física
DCI, Campus León

DIVISIÓN DE CIENCIAS E INGENIERÍAS, CAMPUS LEÓN
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CENTRO DE INVESTIGACIÓN Y DE ESTUDIOS AVANZADOS
DEL INSTITUTO POLITÉCNICO NACIONAL

México, D. F., a 8 de enero 2014

A quien corresponda:

Por medio de la presente les informo que he leído el texto de la tesis "Selected low energy probes for New Physics", que la estudiante Vannia González Macías ha escrito como requisito para la obtención del grado de Doctorado en Física, en la Universidad de Guanajuato.

Considero que este escrito reúne los requisitos de originalidad, claridad, buena documentación e importancia para que sea presentado como Tesis de Doctorado.

Sin otro particular de momento, les envío mis más cordiales saludos.

Atentamente

Dr. Gabriel López Castro
Profesor Investigador
Departamento de Física

Av. Instituto Politécnico Nacional # 2508 Col. San Pedro Zacatenco México, D.F. C.P. 07360
Tel.: 5747-3800 Fax: 5747-7002

Universidad de Guanajuato

León Guanajuato, 07 de Enero de 2014

Dr. Guillermo Mendoza

Director de la División de Ciencias e Ingenierías

Presente

Por este conducto le informo que he leído la tesis titulada "*Selected low energy probes for new physics*", que para obtener el grado de Doctor en Física ha sido formulada por la M. F. Vannia González Macías. Las correcciones sugeridas han sido incorporadas al texto por lo cual tiene mi autorización para proceder a la defensa de la misma.

Sin otro particular, reciba un cordial saludo.

Atentamente

Dr. Mauro Napsuciale Mendívil

Sinodal

c.c.p. archivo

DEPARTAMENTO DE FISICA, DIVISION DE CIENCIAS E INGENIERÍAS, CAMPUS LEÓN

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Universidad de Guanajuato

León, Gto., a 9 de enero de 2014

Dr. Teodoro Cordova Fraga
Coordinador de docencia
División de Ciencias e Ingenierías
Presente

Estimado Dr. Cordova:

Por medio de la presente le informo que he revisado de la tesis de maestría de la **Maestra Vannia González Macías** titulada "Selected Low Energy Probes for New Physics". En base a esta revisión puedo afirmar que los resultados presentados en la tesis permiten que se pueda proceder a su defensa.

Aprovecho la ocasión para enviarle un cordial saludo.

ATENTAMENTE
"LA VERDAD OS HARÁ LIBRES"

DR. FRANCISCO SASTRE CARMONA

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