**Abstract:** We consider the possibility to explain the recent \( R_K \) and \( R_{K^*} \) anomalies in a 2-Higgs Doublet Model, known as Aligned, combined with a low scale seesaw mechanism generating light neutrino masses and mixings. In this class of models, a large Yukawa coupling allows for significant non-universal leptonic contributions, through box diagrams mediated by charged Higgs bosons and right-handed neutrinos, to the \( b \to s\ell^+\ell^- \) transition that can then account for both \( R_K \) and \( R_{K^*} \) anomalies.
1 Introduction

Recently, the LHCb collaboration announced intriguing results [1, 2] for the ratios 
\[ R_K = \frac{\text{BR}(B^+ \to K^+\mu^+\mu^-)/\text{BR}(B^+ \to K^+e^+e^-)} \]
and 
\[ R_{K^*} = \frac{\text{BR}(B^0 \to K^{*0}\mu^+\mu^-)/\text{BR}(B^0 \to K^{*0}e^+e^-)} \].
In fact, it was reported that for two dilepton invariant mass-squared bins, \( R_K \) and \( R_{K^*} \) are given by

\[ R_K = 0.846^{+0.060}_{-0.054}^{+0.016}_{-0.014} \quad \text{for} \quad 1.1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2, \]

\[ R_{K^*} = \begin{cases} 
0.66^{+0.11}_{-0.07} \pm 0.03 \quad & \text{for} \quad 0.045 \text{ GeV}^2 \leq q^2 \leq 1.1 \text{ GeV}^2 \\
0.69^{+0.11}_{-0.07} \pm 0.05 \quad & \text{for} \quad 1.1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2 
\end{cases} \] (1.1)

These measurements contradict the Standard Model (SM) expectations: \( R_K^{SM} \approx R_{K^*}^{SM} \approx 1 \) [3], by a \( \approx 2.5\sigma \) deviation. Hence, they are considered as important hints of new physics that violates lepton universality.

Any signal of possible lepton non-universality would be striking evidence for physics Beyond the SM (BSM). Therefore, the \( R_K \) and \( R_{K^*} \) anomalies attracted the attention of theoretical particle physicists and several new physics scenarios have been proposed to accommodate these results, for a review, see [4] and related references.

A non-trivial flavour structure in the lepton sector is already required, beyond the SM, in order to explain the observation of neutrino oscillations. Therefore, it is plausible and very attractive if the mechanism behind the \( B \) anomalies can be related to the same physics responsible for the non-zero neutrino masses and oscillations. In view of this, one then ought to explore extensions of the SM able to address both phenomena.
One of the possible scenarios for generating the observed lepton non-universality is to allow for large Yukawa couplings of the right-handed neutrinos with Higgs fields and charged leptons, as in a low scale seesaw mechanism, with the inverse seesaw being one of the most notorious examples, for generating light neutrino masses. In this case, a non-trivial neutrino Yukawa matrix may lead to different results for $BR(B \to Ke^+e^-)$ and $BR(B \to K\mu^+\mu^-)$. This framework has been recently considered in the Supersymmetric (SUSY) $B-L$ extension of the SM, where it was shown that the box diagram mediated by a right-handed sneutrino, higgsino-like chargino and light stop can account simultaneously for both $R_K$ and $R_{K^*}$. [5]. In non-SUSY models, a similar box diagram can be obtained through a charged Higgs, instead of a chargino, and a right-handed neutrino, instead of a right-handed sneutrino. Therefore, a rather minimal model that can account for these discrepancies is an extension of the SM with two Higgs doublets (so that we can have a physical charged Higgs boson) and right-handed neutrinos along with a low scale seesaw mechanism (to guarantee large neutrino Yukawa couplings). Note that there have been several attempts at explaining the above results through the penguin diagram as well, with a non-universal $Z'$, and also through tree level mediation of flavour violating $Z'$ or leptoquark that induce a non-universal $b \to s\ell^+\ell^-$ transition. However, these types of flavour violating interactions are subject to severe experimental limits.

The 2-Higgs Doublet Model (2HDM), which is motivated by SUSY and Grand Unified Theories (GUTs), is the simplest model that includes charged Higgs bosons. According to the types of couplings of the two Higgs doublets to the SM fermions doublets and singlets, we may have different type of 2HDMs. For example, if only one Higgs doublet couples to the SM fermions one obtains the type I 2HDM. While, in the case of one Higgs doublet coupling to the up quarks and the second Higgs doublet coupling to the down quarks and charged leptons, one obtains the type II 2HDM. Also, we may have a type III or IV 2HDM if both Higgs doublets couple to both up and down quarks as well as charged leptons. However, severe constraints are imposed on 2HDMs due to their large contributions to Flavour Changing Neutral Currents (FCNCs) that contradict the current experimental limits. Therefore, assumptions on the Yukawa couplings generated by different Higgs doublets are usually imposed. One of these assumptions is the alignment between the Yukawa couplings generated by $\Phi_1$ and $\Phi_2$, the two Higgs doublet fields. This class of models is called the Aligned 2HDM (A2HDM). It is worth mentioning that, as discussed below, other 2HDMs cannot account for the LHCb results of lepton non-universality.

In this paper, we emphasise that in the A2HDM, extended by (heavy) right-handed neutrinos in order to generate the (light) neutrino masses through a seesaw mechanism, interesting results can be obtained for several flavour observables, such as $\mu \to e\gamma$, $B_s \to \mu\mu$ and, indeed, $R_K^{(\ast)}$. In particular, one can account simultaneously for both aforementioned results on $R_K$ and $R_{K^*}$, through a box diagram mediated by a right-handed neutrino, top quark and charged Higgs boson.

The paper is organised as follows. In section 2 we introduce the main features of our A2HDM focusing on the neutrino sector and its interplay with the Higgs structures. (Here, we also briefly review some particular realisations of low scale seesaw models for generating light neutrino masses.) In section 3 we discuss the most relevant constraints
from flavour physics that affect our A2HDM parameter space. The calculation of the A2HDM contributions to $b \rightarrow s\ell^+\ell^-$ transitions mediated by charged Higgs bosons and right-handed neutrinos is given in section 4. Our numerical results are presented in section 5. Finally, our conclusions and remarks are given in section 6.

2 The A2HDM

The 2HDM is characterised by two Higgs doublets with hypercharge $Y = 1/2$ which, in the Higgs basis, can be parameterised as

$$
\Phi_1 = \left( \frac{1}{\sqrt{2}} (v + \phi_1^0 + iG^0) \right), \quad \Phi_2 = \left( \frac{1}{\sqrt{2}} (\phi_2^0 + i\phi_3^0) \right),
$$

(2.1)

where $v$ is the Electro-Weak (EW) Vacuum Expectation Value (VEV) and $G^+, G^0$ denote the Goldstone bosons. The two doublets describe five physical scalar degrees of freedom which are given by the two components of the charged Higgs $H^\pm$ and three neutral states $\phi_0^i = \{h, H, A\}$, the latter obtained from the rotation of the $\phi_0^i$ fields into the mass eigenstate basis. The scalar squared mass matrix $M^2_S$ is determined by the structure of the 2HDM scalar potential, see for instance [6–8], and diagonalised by the orthogonal matrix $R$, where

$$
R M^2_S R^T = \text{diag}(M^2_h, M^2_H, M^2_A), \quad \phi_0^i = R_{ij} \phi_0^j.
$$

(2.2)

In general, the three mass eigenstates $\phi_0^i$ do not have definite CP transformation properties but in the CP-conserving scenario $\phi_0^3$ does not mix with the other two neutral states and the scalar spectrum consists of a CP-odd field $A = \phi_0^3$ and two CP-even fields $h$ and $H$ that are defined from the interaction eigenstates through the two-dimensional orthogonal matrix

$$
\begin{pmatrix}
  h \\
  H
\end{pmatrix} = 
\begin{pmatrix}
  \cos \alpha & \sin \alpha \\
  -\sin \alpha & \cos \alpha
\end{pmatrix}
\begin{pmatrix}
  \phi_0^1 \\
  \phi_0^2
\end{pmatrix}.
$$

(2.3)

The most general Yukawa Lagrangian of the 2HDM can be written in the Higgs basis as

$$
-\mathcal{L}_Y = \bar{Q}'_L (Y'_{1d} \Phi_1 + Y'_{2d} \Phi_2) d'_R + \bar{Q}'_L (Y'_{1d} \Phi_1 + Y'_{2d} \Phi_2) u'_R \\
+ \bar{L}'_L (Y'_{1u} \Phi_1 + Y'_{2u} \Phi_2) \ell'_R + \bar{L}'_L (Y'_{1u} \Phi_1 + Y'_{2u} \Phi_2) \nu'_R + \text{h.c.}
$$

(2.4)

where the quark $Q'_L, u'_R, d'_R$ and lepton $L'_R, \ell'_R, \nu'_R$ fields are defined in the weak interaction basis and we also included the couplings of the left-handed lepton doublets with the right-handed neutrinos. The $\Phi_{1,2}$ fields are the two Higgs doublets in the Higgs basis and, as customary, $\Phi_i = i\sigma^2 \Phi_i^*$. The Yukawa couplings $Y'_{1j}$ and $Y'_{2j}$, with $j = u, d, \ell$, are $3 \times 3$ complex matrices while $Y'_{1u}$ and $Y'_{2u}$ are $3 \times n_R$ matrices, with $n_R$ being the number of right-handed neutrinos. In general, the Yukawas $Y'_1$ and $Y'_2$ cannot be simultaneously diagonalised in flavour space, so that, whilst the quark and the charged-lepton $Y'_1$ can be recast into a diagonal form in the fermion mass eigenstate basis, namely, $Y_1 = \sqrt{2}/vM$, with $M$ being
the fermion mass matrix, $Y_2$ would remain non-diagonal and thus give rise to potentially dangerous tree-level FCNCs. This problem is usually solved by enforcing that only one of the two Higgs doublets couple to a given right-handed field. This requirement is satisfied by implementing a discrete $Z_2$ symmetry acting on the Higgs and fermion fields. There are four non-equivalent choices: type I, II, III and IV (as previously intimated). Another general way to avoid tree-level FCNCs in the Higgs sector is to require the alignment, in flavour space, of the two Yukawa matrices that couple to the same right-handed fermion [9], namely,

$$Y_{2,d} = \zeta_d Y_{1,d} \equiv \zeta_d Y_d, \quad Y_{2,u} = \zeta_u^* Y_{1,u} \equiv \zeta_u^* Y_u, \quad Y_{2,\ell} = \zeta_\ell Y_{1,\ell} \equiv \zeta_\ell Y_\ell,$$

(2.5)

where the proportionality constants $\zeta_f$ are arbitrary family universal complex parameters. This scenario is dubbed A2HDM. The allowed sources of FCNCs at quantum level are highly constrained and the resulting structures are functions of the mass matrices and Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, so that this model provides an explicit example of the popular Minimal Flavour Violation (MFV) scenario [10].

Even though the alignment of the Yukawa matrices is strictly required, from observations, only in the quark and charged lepton sectors, we assume that the same mechanism which guarantees the aligned structure in the SM flavour space also holds in the neutrino sector and leads to

$$Y_{2,\nu} = \zeta_\nu^* Y_{1,\nu} \equiv \zeta_\nu^* Y_\nu.$$

(2.6)

In all sectors, the alignment is fixed to be exact at some specified scale $\mu_0$ and subsequently will misalign due to radiative corrections, as discussed in [11, 12]. However, the flavour structure of the model constrain the nature of the new sourced of FCNCs induced by Renormalisation Group effects. Quantitatively, in the quark sector, these FCNC contributions are suppressed by mass hierarchies $m_q m_q^2/v^3$ and provide negligible effects [11, 12]. We will not consider the impact of the misalignment in this work.

Interestingly, $\zeta_f$ can provide new sources of CP violation but in this work we will consider only real values. Notice also that the usual 2HDMs in which tree-level FCNCs are removed by exploiting the discussed $Z_2$ discrete symmetry, namely the type I, II, III and IV, can be recovered for particular values of the proportionality constants $\zeta_f$ as shown in Tab. 1.

| Table 1. Relation between the $\zeta_f$ couplings of the A2HDM and the ones of the $Z_2$ symmetric scenarios. |
|---|---|---|---|---|
| $\zeta_u$ | $\cot\beta$ | $\cot\beta$ | $\cot\beta$ | $\cot\beta$ |
| $\zeta_d$ | $\cot\beta$ | $-\tan\beta$ | $-\tan\beta$ | $\cot\beta$ |
| $\zeta_\ell$ | $\cot\beta$ | $-\tan\beta$ | $\cot\beta$ | $-\tan\beta$ |

The Yukawa Lagrangian in Eq. (2.4) generates a Dirac mass matrix for the standard neutrinos and can also be supplemented by a Majorana mass term $M'_R$ for the right-handed...
Due to the alignment of the Yukawa matrices all the couplings of the scalar fields to fermions where the couplings of the neutral Higgs states to the fermions are given by

\[ U_L^\dag Y \nu U_R^\nu = Y \nu \equiv \sqrt{2} \text{diag}(m_e, m_\mu, m_\tau), \]
\[ U_R^\nu T M_R^\nu U_R^\nu = M_R \equiv \text{diag}(M_1, \ldots, M_{n_R}), \]

while \( Y_\nu = U_L^\dag Y \nu U_R^\nu \) remains non-diagonal. In this basis the neutrino mass matrix can be written as

\[-\mathcal{L}_{M_\nu} = \frac{1}{2} N_L^T C M N_L + \text{h.c.} = \frac{1}{2} (\nu_L^T \nu_R^\nu)^T C \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix} \begin{pmatrix} \nu_L \\ \nu_R^\nu \end{pmatrix}, \]

with \( M_D = \frac{\sqrt{2}}{v} Y_\nu^\nu \) being the neutrino Dirac mass. This can be diagonalised with the unitary \((3 + n_R) \times (3 + n_R)\) matrix \( U \),

\[ \begin{pmatrix} \nu_L \\ \nu_R^\nu \end{pmatrix} = U \begin{pmatrix} \nu_l \\ \nu_h \end{pmatrix} \equiv \begin{pmatrix} U_{Li} & U_{Lh} \\ U_{Ri}^\dag & U_{Rh}^\dag \end{pmatrix} \begin{pmatrix} \nu_l \\ \nu_h \end{pmatrix}, \]

such that \( M_\nu = U^T M U \) provides the masses of the three light active neutrinos \( \nu_l \) and of the remaining \( n_R \) heavy sterile neutrinos \( \nu_h \).

The Yukawa interactions of the physical scalars with the mass eigenstate fermions are then described by

\[-\mathcal{L}_Y = \frac{\sqrt{2}}{v} \left[ \bar{u}(-\zeta_u m_u V_{ud} P_L + \zeta_d V_{ud} m_d P_R) d + \bar{\nu}_l(-\zeta_\nu m_\nu U_L^\dag P_L + \zeta_\ell U_L^\dag m_\ell P_R) \ell \\
+ \bar{\nu}_h(-\zeta_\nu m_\nu U_L^\dag P_L + \zeta_\ell U_L^\dag m_\ell P_R) \ell \right] H^+ + \text{h.c.} \]
\[ + \frac{1}{v} \sum_i \sum_{f=u,d,e} \xi_{i,f}^0 \bar{f} m_f P_R f \\
+ \frac{1}{v} \sum_i \xi_{i,\ell} (\bar{\nu}_l U_L^\dag + \bar{\nu}_h U_L^\dag) P_R (U_L m_\nu \nu_l + U_L m_\nu \nu_h) + \text{h.c.} \]

where the couplings of the neutral Higgs states to the fermions are given by

\[ \xi_{u,\nu} = \mathcal{R}_{i1} + (\mathcal{R}_{i2} - i\mathcal{R}_{i3}) \zeta_u^*, \quad \xi_{d,\ell} = \mathcal{R}_{i1} + (\mathcal{R}_{i2} + i\mathcal{R}_{i3}) \zeta_d, \ell. \]

Due to the alignment of the Yukawa matrices all the couplings of the scalar fields to fermions are proportional to the corresponding mass matrices. Finally, the weak neutral and charged interactions of the neutrinos are

\[ \mathcal{L}_Z = \frac{g}{2 \cos \theta_W} (\bar{\nu}_l U_L^\dag + \bar{\nu}_h U_L^\dag) \gamma^\mu (U_L \nu_l + U_L \nu_h) Z_\mu, \]
\[ \mathcal{L}_W = -\frac{g}{\sqrt{2}} \left[ (\bar{\nu}_l U_L^\dag + \bar{\nu}_h U_L^\dag) \gamma^\mu P_R \ell \right] W_\mu^\nu + \text{h.c.} \]
In this paper, rather than presenting a complete model in the neutrino sector by specifying the structure and the hierarchies of the neutrino mass matrices, we work in a simplified framework that captures the interesting phenomenology whilst preserving a significant degree of model independence. In particular, we consider a single extra heavy neutrino despite the usual requirement of additional sterile states to fully accommodate the observed pattern of the light neutrino masses and mixing angles. Indeed, low scale right-handed neutrinos with sizeable mixings with the SM left-handed neutrino states, such that they may provide visible effects in physical observables at the EW scale, usually affect the light neutrino masses with inadmissible large contributions. This issue is nicely solved in extended seesaw models [13–20], as for example the linear or inverse seesaw mechanism, in which extra sterile neutrino states are introduced to allow for large mixings while correctly reproducing the observed smallness of the light neutrino masses.

For the sake of definiteness, in the following section we briefly present some specific setup that could be employed to realise the phenomenological scenario described above.

2.1 Some explicit examples of low scale seesaw mechanism

As stated above, in order to achieve a sizeable mixing with the heavy sterile neutrino while avoiding, at the same time, large contributions to the light neutrino masses, it is necessary to require the neutrino Majorana mass matrix and the Dirac neutrino Yukawa coupling to realise a particular structure. In order to match the nomenclature usually employed in the literature, we split the set of right-handed neutrinos defined above into two classes, one of $n_S$ SM-singlet fermionic fields $S^i$ and another one of fermionic states that we continue to call right-handed neutrinos. The differences between the two will be clear in a moment.

In the basis $N_L = (\nu_L, \nu_R, S)^T$, the mass matrix can be parameterised as

$$
\mathcal{M} = \begin{pmatrix}
0 & m_D & m_S \\
m_D^T & m_N & m_R \\
m_S^T & m_R^T & \mu_S
\end{pmatrix}
$$

(2.14)

which has the same structure of the one given in Eq. (2.9) provided that

$$
M_D \equiv (m_D, m_S), \quad M_R \equiv \begin{pmatrix}
m_N & m_R \\
m_R^T & \mu_S
\end{pmatrix}.
$$

(2.15)

The $m_D$ and $m_S$ are, respectively, $3 \times n_R$ and $3 \times n_S$ mass matrices mediating the interactions between the charged leptons and the right-handed and sterile neutrinos while $m_N$, $m_R$ and $\mu_S$ are $n_R \times n_R$, $n_R \times n_S$ and $n_S \times n_S$ mass matrices, respectively. As both right-handed $\nu_R$ and sterile $S$ neutrinos can be assigned lepton number $L = 1$, the mass terms $m_N$, $m_S$ and $\mu_S$ violate lepton number by two units.

Two commonly studied mass patterns are the inverse and linear seesaws, which are characterised by $m_S = m_N = 0$ and $\mu_S = m_N = 0$, respectively. In these cases, the vanishing of the Majorana mass $\mu_S$ or $m_S$ would restore lepton number conservation and, as such, would increase the symmetry of the model. This feature makes the two masses naturally small accordingly to ’t Hooft naturalness principle.
Following the standard seesaw calculation and by assuming the hierarchy $\mu_S(m_S) \ll m_D, m_R$ for the inverse (linear) seesaw scenario, the $3 \times 3$ light neutrino mass matrix is

$$m_{\text{light}} \simeq \begin{cases} m_D(m_T^R)^{-1} \mu_S m_R^{-1} m_D^T & \text{inverse seesaw} \\ m_S m_R^{-1} m_D^T + m_D(m_T^R)^{-1} m_S & \text{linear seesaw} \end{cases}$$

which is diagonalised by the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix $U_{\text{PMNS}}$, namely

$$U_{\text{PMNS}}^T m_{\text{light}} U_{\text{PMNS}} = m_\nu \equiv \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}).$$

Differently from the standard type I seesaw case, in which $m_\nu \sim m_D^2/m_R$ with $m_D \ll m_R$, the lightness of the active neutrino masses is ensured in these low scale seesaw scenarios by the smallness of the $\mu_S(m_S)$ parameters. This feature prevents $m_D$ to be extremely suppressed with respect to the Majorana mass and, as such, may allow for non-negligible couplings between the heavy neutrinos and the SM gauge bosons which are set by the mixing $U_{Lh}$.

In order to understand the dependence of the mixing $U_{Lh}$ of the left-handed SM neutrinos with the extra sterile states, it is instructive to study the mixing matrix $U$ in the limit of negligible $\mu_S(m_S)$ and small $m_D/m_R$. While the first requirement is necessary to reproduce the lightness of the active neutrino states, the latter is used here only to simplify the structure of $U$ which reads as

$$U = \begin{pmatrix} 1 & \frac{1}{\sqrt{2}} m_D^* m_R^{-1} & \frac{i}{\sqrt{2}} m_D^* m_R^{-1} \\ 0 & \frac{1}{\sqrt{2}} & \frac{-i}{\sqrt{2}} \\ -m_R^{-1} m_D^T & \frac{i}{\sqrt{2}} & \frac{1}{\sqrt{2}} \end{pmatrix} + \mathcal{O}\left(\frac{m_D^2}{m_R^2}\right).$$

Notice that the PMNS matrix has been set to the unit one, consistently with the approximation $m_{\nu_i} \simeq 0$. From Eq. (2.18), one can immediately realise that $U_{Lh}$ is set, as naively expected from dimensional arguments, by the ratio $m_D/m_R$.

Once some specific inputs are provided for $m_D$ and $m_R$, the corresponding $\mu_S(m_S)$ matrix that ensures the agreement with the light neutrino mass splittings and mixing angles can always be reconstructed from Eq. (2.16). For instance, in the inverse seesaw case, we find

$$\mu_S = m_R^T m_D^{-1} U_{\text{PMNS}}^* m_\nu U_{\text{PMNS}}^T (m_D^T)^{-1} m_R,$$

where $m_\nu$ and $U_{\text{PMNS}}$ are chosen in agreement with the bounds from the low-energy neutrino data which we report below for the sake of completeness. In particular, one should enforce the following constraints from the latest results of the $\nu$fit group [21] extracted from the $\nu$fit 3.2 (2018) data.

1) **Neutrino mass squared differences**  
The 3$\sigma$ Confidence Level (CL) ranges on the mass squared differences

$$\Delta m_{21}^2 = (6.80 \rightarrow 8.02) \times 10^{-5}\text{eV}^2$$

$$\Delta m_{3l}^2 = \begin{cases} (2.399 \rightarrow 2.593) \times 10^{-3}\text{eV}^2 & \text{(for } \ell = 1 \text{ N.O.}) \\ (-2.562 \rightarrow -2.369) \times 10^{-3}\text{eV}^2 & \text{(for } \ell = 2 \text{ I.O.)} \end{cases}$$

\(2.20\)
where the first and second possibility refer to the assumption of normal and inverted ordering in the light neutrino masses, respectively.

2) Leptonic mixing matrix

The $3\sigma$ CL ranges on the elements of the leptonic mixing matrix $U_{\text{PMNS}}$

$$\sin^2 \theta_{12} = (0.272 \rightarrow 0.346)$$
$$\sin^2 \theta_{23} = \begin{cases} (0.418 \rightarrow 0.613) & \text{N.O.} \\ (0.435 \rightarrow 0.616) & \text{I.O.} \end{cases}$$
$$\sin^2 \theta_{13} = \begin{cases} (0.01981 \rightarrow 0.02436) & \text{N.O.} \\ (0.02006 \rightarrow 0.02452) & \text{I.O.} \end{cases} \quad (2.21)$$

Notice also that $m_D$ and $m_R$ cannot be chosen freely since the $U_{Lh}$ block of the mixing matrix that they define is constrained by the unitarity requirement that directly affects the analysis presented in this work. The corresponding bound is discussed in the next section together with all the other relevant constraints.

3 Relevant parameter space and constraints

3.1 Unitarity bounds on the neutrino mixing matrix

The $3 \times 3$ block of the mixing matrix $U$ corresponds to a non-unitary $\tilde{U}_{\text{PMNS}}$ matrix. The bounds on the deviation from unitarity of $\tilde{U}_{\text{PMNS}}$ have been obtained in [22, 23] using an effective field theory approach in which the masses of the heavy neutrinos lie above the EW scale. This constraint can be recast as follows

$$\epsilon_{\alpha\beta} \equiv \left| \sum_i^n U_{\alpha i}^* U_{\beta i} \right| \equiv \left| \sum_i^n (U_{Lh}^*)_{\alpha i} (U_{Lh})_{\beta i} \right| = \left| \delta_{\alpha\beta} - (\tilde{U}_{\text{PMNS}}^* \tilde{U}_{\text{PMNS}})_{\alpha\beta} \right|.$$\quad (3.1)

$$\left| \tilde{U}_{\text{PMNS}} \tilde{U}_{\text{PMNS}}^\dagger \right| = \begin{pmatrix} (0.9979 \rightarrow 0.9998) & < 10^{-5} & < 0.0021 \\ < 10^{-5} & (0.9996 \rightarrow 1.0) & < 0.0008 \\ < 0.0021 & < 0.0008 & (0.9947 \rightarrow 1.0) \end{pmatrix}.\quad (3.1)$$

3.2 Lepton flavour violating processes

We consider the lepton flavour violating decays $\ell_\alpha \to \ell_\beta \gamma$ induced at one-loop order by the sterile neutrinos and check their compatibility with the experimental upper bounds at 90% CL [24],

$$\text{BR}(\mu \to e\gamma) \leq 4.2 \times 10^{-13}, \quad \text{BR}(\tau \to e\gamma) \leq 3.3 \times 10^{-8}, \quad \text{BR}(\tau \to \mu\gamma) \leq 4.4 \times 10^{-8}. \quad (3.2)$$

The Branching Ratios (BRs) of the aforementioned decay rates are given by

$$\text{BR}(\ell_\alpha \to \ell_\beta \gamma) = C \left| \sum_i^n (U_{Lh})_{\alpha i} (U_{Lh})_{\beta i} \left[ \mathcal{G}_W \left( \frac{m_{\nu_{\beta i}}}{M_W^2} \right) + \mathcal{G}_H \left( \frac{m_{\nu_{\beta i}}}{M_W^2} \right) \right] \right|^2 \quad (3.3)$$

with

$$C = \frac{\alpha_W^3 s_W^2}{256 \pi^2} \left( \frac{m_{\ell_\alpha}}{M_W} \right)^4 \frac{m_{\ell_\alpha}}{\Gamma_{\ell_\alpha}} \quad (3.4)$$
where $\Gamma_{\ell^\alpha}$ is the total decay width of the lepton $\ell^\alpha$ and the loop functions are
\begin{align*}
G_{W^\pm}(x) &= \frac{-x + 6x^2 - 3x^3 - 2x^4 + 6x^3 \log x}{4(x - 1)^4}, \\
G_{H^\pm}(x) &= \frac{\zeta_\nu^2}{3} G_{W^\pm}(x) + \zeta_\nu \zeta_\ell \frac{x(-1 + x^2 - 2x \log x)}{2(x - 1)^3},
\end{align*}
(3.5)
where we have neglected the mass of the lepton in the final state.

The $G_{H^\pm}$ can offer large contributions, larger than $G_{W^\pm}$, for sizeable values of the couplings $\zeta_\nu, \zeta_\ell$, which are, as such, strongly constrained by lepton flavour violating processes. These can be tamed by controlling the size of the mixing matrix elements $(U_{Lh})_{\alpha i}$ which should be highly suppressed in order to avoid any large contribution to these sensitive processes.

### 3.3 Flavour constraints from meson processes

Here we briefly mention the relevant constraints from measurement of flavour observables in meson mixing and decays. These have been studied in the context of general 2HDMs and the majority of the bounds extracted from these can be straightforwardly applied in our case since the presence of the sterile neutrinos does not add any significant contribution at leading order. In particular, the 2HDM with the alignment in the flavour sector has been scrutinised in [11, 12, 25, 26] to which we refer for delineating the allowed parameter space spanned by the $\zeta_{u,d,\ell}$ couplings and the charged Higgs mass $M_{H^\pm}$.

- **Neutral meson mixing**
  The meson mixing observables $\Delta M_s$, $\Delta M_d$ and $|\epsilon_K|$ constrain large value of the $\zeta_u$ parameter. In turn, they do not significantly affect $\zeta_d$ whose dependence is suppressed by $m_b^2/M_W^2$ in the first two observables while does not appear at all in the third one. Obviously, there is no dependence on $\zeta_\ell$ and $\zeta_\nu$.

- **Radiative $B_s \to X_s \gamma$ decay**
  The $B_s \to X_s \gamma$ decay rate represents one of the best measured observables and it is employed to constrain several new physics scenarios. The contribution of the charged Higgs boson is encoded in the Wilson coefficients $C_7$ and $C_8$. At leading order, the corresponding new physics corrections are sensitive to $\zeta_u$ and $\zeta_d$ and are given by
  \begin{align*}
  C_i &= \frac{(\zeta_u)^2 G^1_i}{3} \left( \frac{M_t^2}{M_{H^\pm}^2} \right) + \zeta_u \zeta_d G^2_i \left( \frac{M_t^2}{M_{H^\pm}^2} \right),
  \end{align*}
  (3.6)
  with
  \begin{align*}
  G^1_7(x) &= \frac{y(7 - 5y - 8y^2)}{24(y - 1)^3} + \frac{y^2(3y - 2)}{4(y - 1)^4} \log x, \\
  G^2_7(x) &= \frac{y}{12(y - 1)^2} + \frac{y(3y - 2)}{6(y - 1)^3} \log x, \\
  G^1_8(x) &= -\frac{3y^2}{8(y - 1)^3} - \frac{3y^2}{4(y - 1)^4} \log x, \\
  G^2_8(x) &= \frac{y(3 - y)}{4(y - 1)^2} - \frac{y}{2(y - 1)^3} \log x.
  \end{align*}
  (3.7)
  As for the new physics contributions to the meson mixing observables, the values of the $\zeta_\ell$ and $\zeta_\nu$ parameters are completely irrelevant in the determination of the $b \to s \gamma$ transition.
Leptonic decay of the neutral mesons $B^0_q \rightarrow \mu^+ \mu^-$

The BR of the $B^0_q \rightarrow \mu^+ \mu^-$ meson decay is given by

$$
\text{BR}(B^0_q \rightarrow \mu^+ \mu^-) = \text{BR}_{\text{SM}}(B^0_q \rightarrow \mu^+ \mu^-)(|P|^2 + |S|^2)
$$

with

$$
P = \frac{C_{10}}{C_{10}^{\text{SM}}} + \frac{m_{B^0_q}^2}{2M_W^2} \left( \frac{m_b}{m_b + m_q} \right) C_P, \quad S = \sqrt{1 - \frac{4m_{\mu}^2}{m_{B^0_q}^2} \frac{m_{B^0_q}^2}{2M_W^2} \left( \frac{m_b}{m_b + m_q} \right) \frac{C_S}{C_{10}^{\text{SM}}}}. \quad (3.9)
$$

The leading new physics contribution appears in the Wilson coefficient $C_{10}$ and it is mediated by the charged Higgs in 2HDMs. Interestingly, the same coefficient is also corrected by the presence of the heavy sterile neutrinos and it is sensitive to, besides $\zeta_{u,d,\ell}$, the coupling $\zeta_{\nu}$ of the charged Higgs to the sterile neutrino states. The impact of the $C_F$ and $C_S$ coefficients is, in contrast, suppressed by the $m_{B^0_q}^2/M_W^2$ factor unless the former can be enhanced with respect to $C_{10}^{\text{SM}}$ as, for instance, in a $Z_2$ symmetric model with large $\tan \beta$. In the flavour aligned 2HDM and in the parameter space in which we are interested in, namely $\zeta_u \simeq 1$ and $\zeta_d \simeq \zeta_{\ell} \simeq 0$, the bound from the measurement of the $B^0_s \rightarrow \mu^+ \mu^-$ transition is usually weaker than the one from the $B_s \rightarrow X_s \gamma$ decay. Nevertheless, due to the contribution from the heavy sterile neutrinos which is proportional to the new parameter $\zeta_{\nu}$, we recalculate the corresponding constraint by employing the flavio package [27].

Leptonic decay of the charged mesons $M^\pm \rightarrow \tau^\pm \nu$

The $M^\pm \rightarrow \tau^\pm \nu$ decay occurs at tree level through charged current processes and the corresponding BR is

$$
\text{BR}(M^\pm \rightarrow \tau^\pm \nu) = \frac{\tau_M G_F^2 m_M m_{\tau_M}^2}{8\pi} \left( 1 - \frac{m_{\tau_M}^2}{m_M^2} \right)^2 |V_{ud}|^2 f_M^2 |1 + C_H|^2, \quad (3.10)
$$

where $f_M$ and $\tau_M$ are the decay constant and the lifetime, respectively, and the contribution of the charged Higgs boson is encoded in

$$
C_H = \frac{\zeta_u \zeta_{\ell} m_u - \zeta_d \zeta_{\ell} m_d}{m_u + m_d} \frac{m_{\tau_M}^2}{M_{H^\pm}^2}.
$$

(3.11)

The most constraining decay mode is found to be $B \rightarrow \tau \nu$ which is, however, only relevant for light charge Higgs masses.
4 Contributions to $b \to s\ell^+\ell^-$ processes

The effective Hamiltonian for the $b \to s\ell^+\ell^-$ transitions is given by

$$\mathcal{H}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} V_{ts} V_{tb} \sum_i \Delta C_i O_i + \text{h.c.} \tag{4.1}$$

where the relevant operators for the analysis of the $R_K$ and $R_{K^*}$ anomalies are

$$O_9 = (\bar{b} \gamma^\mu P_L s)(\bar{\ell} \gamma_\mu \ell), \quad O_{10} = (\bar{b} \gamma^\mu P_L s)(\bar{\ell} \gamma_\mu \gamma_5 \ell). \tag{4.2}$$

The new physics effects in the corresponding Wilson coefficients can be recasted as

$$\Delta C_i = C_i^{(\Delta_Z)}(H^\pm) + C_i^{(\Delta_Z)}(N_R) \quad \text{Z Penguin}$$

$$+ C_i^{(\Delta_\gamma)}(H^\pm) + C_i^{(\Delta_\gamma)}(N_R) \quad \text{\gamma Penguin}$$

$$+ C_i^{(\Box)}(H^\pm) + C_i^{(\Box)}(N_R) + C_i^{(\Box)}(N_R, H^\pm) \quad \text{Box} \tag{4.3}$$

where $\Delta_Z$, $\Delta_\gamma$ and $\Box$ denote the contributions from the Z penguins, the photon penguins and the box diagrams, respectively. Moreover, $C_i^{(*)}(H^\pm)$ represents the charged Higgs contribution typical of the 2HDM, $C_i^{(*)}(N_R)$ denotes the loop corrections from heavy sterile neutrinos and $W^\pm$ bosons that is present in the seesaw extensions of the SM without extra Higgses while $C_i^{(*)}(N_R, H^\pm)$ represents a combined contribution from diagrams with both sterile neutrinos and charged Higgs.

As the sterile neutrinos are not charged under the (colour) $SU(3)$ gauge group, their contribution to the penguin diagrams is identically zero at leading order, namely, $C_i^{(\Delta_Z)}(N_R) = C_i^{(\Delta_\gamma)}(N_R) = 0$. Moreover, $C_i^{(\Box)}(H^\pm) = C_i^{(\Delta_\gamma)}(H^\pm) = 0$ since in the 2HDM the charged Higgs contributes only to the Z penguin diagram. Finally, $C_i^{(\Box)}(N_R, H^\pm)$ includes the contributions of the heavy neutrinos exchange mediated by a charged Higgs current and, as such, is peculiar of the model considered in this paper. The analytic expressions of the new physics contributions to the Wilson coefficients for the penguin diagrams are as follows.

Penguins: contributions from the charged Higgs boson

$$C_9^{(\Delta_Z)}(H^\pm) = -c_u^2 \left( \frac{1 + 4s_w^2}{8s_w^2} \right) \frac{x_H x_1 (x_H - 1 - \log x_H)}{(x_H - 1)^2}, \tag{4.4}$$

$$C_{10}^{(\Delta_Z)}(H^\pm) = -c_u^2 \left( \frac{1}{8s_w^2} \right) \frac{x_H x_1 (x_H - 1 - \log x_H)}{(x_H - 1)^2}, \tag{4.4}$$

$$C_9^{(\Delta_\gamma)}(H^\pm) = c_u^2 \frac{x_H}{108(x_H - 1)^4} \left[ \frac{1}{6} \left( 3x_H^3 - 6x_H + 4 \right) \log x_H - (x_H - 1)(x_H(47x_H - 79) + 38) \right], \tag{4.4}$$

$$C_{10}^{(\Delta_\gamma)}(H^\pm) = 0. \tag{4.4}$$

Penguins: contributions from the heavy neutrinos

$$C_i^{(\Delta_Z)}(N_R) = C_i^{(\Delta_\gamma)}(N_R) = 0. \tag{4.5}$$
The analytic expressions of the new physics contributions to the Wilson coefficients for the box diagrams are as follows:

**Box: contributions from the charged Higgs boson**

\[ C_9^{(\Box)}(H^\pm) = C_{10}^{(\Box)}(H^\pm) = 0. \]  

**Box: contributions from the heavy neutrinos**

\[ C_9^{(\Box)}(N_R) = \sum_{i=1}^{n_R} |(U_{Lh})_{li}|^2 \frac{x_t}{16 s_W^2 (x_{N_i} - 1)(x_t - 1)^2 (x_{N_i} - x_t)^2} \left[ (x_{N_i} - 1)((x_t - 1)(x_t - x_{N_i}) \right. \\
\left. \times (7x_t - 4x_{N_i}) - (4x_{N_i}^2 - x_{N_i}x_t(3x_t + 8) + x_t^2(6x_t + 1)) \log x_t \right) \\
\left. - (x_t - 1)^2 (4x_{N_i}^2 - 8x_{N_i}x_t + x_t^2) \log x_{N_i} \right], \]

\[ C_{10}^{(\Box)}(N_R) = -C_9^{(\Box)}(N_R). \]  

**Box: contributions from the charged Higgs boson and the heavy neutrinos**

\[ C_9^{(\Box)}(N_R, H^\pm) = \sum_{i=1}^{n_R} |(U_{Lh})_{li}|^2 \left\{ \frac{\zeta_{2}^{\nu}}{16 s_W^2 (x_H - 1)^2 (x_{N_i} - 1)(x_H - x_{N_i})} \right. \\
\left. \times \left[ -x_H(x_H - 1)^2 \log x_{N_i} - (x_{N_i} - 1)((x_H - 1)(x_H - x_{N_i}) \\
+ x_H(-2x_H + x_{N_i} + 1) \log x_H \right] \right. \\
\left. + \frac{\zeta_{\nu}}{8 s_W^2 (x_H - 1)(x_{N_i} - 1)(x_t - 1)(x_H - x_{N_i})(x_H - x_t)(x_{N_i} - x_t)} \right. \\
\left. \times \left[ (x_{N_i} - 1)(x_t - 1)(4x_H - x_t) \log x_H(x_{N_i} - x_t) \\
+ (x_H - 1)((x_t - 1)(x_H - x_t)(x_t - 4x_{N_i}) \log x_{N_i} + 3(x_{N_i} - 1)x_t(x_H - x_{N_i}) \log x_t) \right] \right\}, \]

\[ C_{10}^{(\Box)}(N_R, H^\pm) = -C_9^{(\Box)}(N_R, H^\pm). \]  

In the previous equations we used the following mass ratios

\[ x_t = \frac{M_t^2}{M_W^2}, \quad x_H = \frac{M_H^2}{M_{H^\pm}^2}, \quad x_{N_i} = \frac{M_{N_i}^2}{m_{\ell_{hi}}^2}. \]  

We now discuss some interesting features concerning the structure of these coefficients. Firstly, one can see for the box diagram contribution that the heavy neutrinos give \( C_9 = -C_{10} \). One can also see that the SM box contribution with the light neutrinos, not shown above, is rescaled by \( \sum_{i=1}^{3} |\bar{U}_{\ell}\bar{U}_{PMNS}|^2 \simeq 1 - \eta^2 \), where \( \ell = e, \mu \) and \( \eta^2 \) controlling the departure from unitarity of the PMNS matrix. (Since \( \eta^2 \) is expected to be small and no new physics enhancement factors are present in the SM diagram, we can safely neglect this correction.) Concerning the box diagrams, there is also a new contribution \( C_i^{(\Box)}(N_R) \)
with the heavy neutrinos and virtual $W^\pm$s. This is proportional to $\sum_{i=1}^{n_R} |(U_{Lh})_{li}|^2$ and we do not expect, as confirmed by the numerical analysis, that this contribution can provide large effects onto flavour observables. Indeed, the coupling of the heavy neutrinos to the leptons mediated by the $W^\pm$ boson is fixed by the gauge invariance and proportional to the $SU(2)$ weak gauge coupling. In order to allow for more freedom one has to rely on an extra charged degree of freedom with the simplest possibility being the charged scalar of a 2HDM extension. These contributions are encoded into $C^{(\square)}_i(N_{R}, H^\pm)$. The $Z_2$ symmetric scenarios of the 2HDM are among the simplest ones but barely produce significant effects in the $C_9,10$ Wilson coefficients. This can be understood from Tab. 1, because the corrections to $C_{9,10}$ would be proportional to $\zeta^2_\nu = \zeta^2_\ell = \cot^2 \beta$, independently from the specific realisation, and thus relevant only for $\tan \beta < 1$, which is severely constrained by $b \to s\gamma$. The A2HDM allows to disentangle $\zeta_\alpha$ from $\zeta_\nu$, such that, while the former is still bound from $b \to s\gamma$, the latter can be varied freely thus providing significant contributions to the Wilson coefficients in some region of the parameter space. We recall again that the alignment in the neutrino sector is not strictly required by the flavour physics but we, nevertheless, impose it by assuming that the same mechanism ensuring the proportionality between the Yukawa couplings is in place in both the quark and lepton sectors.

5 Results

- 13 –
Among the different lepton flavour violating processes, $\mu \rightarrow e\gamma$ is the most constraining one. In Fig. 2 we show the corresponding BR generated by a single heavy neutrino as a function of the squared mixing angle of the same heavy neutrino with the electron one. The other parameters have been fixed as $M_{H^\pm} = 700$ GeV, $m_{\nu_{hi}} = 500$ GeV, the $\zeta_\ell = 0$ and $|\langle U_{Lh}\rangle_{i\mu}|^2 = 0.4 \times 10^{-3}$. The latter corresponds to the maximum allowed value from unitarity constraints, see Eq. (3.1). Since in a model with a single Higgs doublet we would obtain $\text{BR}/U^4 \sim (6 - 7) \times 10^{-4}$ in the heavy neutrino mass range of $500 - 1000$ GeV, it is clear that the diagrams with the charged Higgs give the dominant and a very large contribution to the BR. Using the largest values for the mixing angles allowed by unitarity we obtain a suppression of the BR of $\sim 10^{-7}$ which cannot accommodate the strong bound on $\mu \rightarrow e\gamma$ from the MEG experiment. This suggest that, if a low scale seesaw is embedded into a 2HDM framework, a given heavy neutrino (or, in a scenario with a large hierarchy between heavy neutrino states, the lightest one) may have a non-negligible mixing only with SM neutrinos of a given flavour eigenstate otherwise the charged Higgs boson would induce unacceptably large effects on lepton flavour violating processes. The two realisations with SM neutrinos of a given flavour eigenstate otherwise the charged Higgs boson would would be $\langle U_{Lh}\rangle_{i\mu} \not\equiv 0, \langle U_{Lh}\rangle_{ie} \simeq 0$ or $\langle U_{Lh}\rangle_{i\mu} \simeq 0, \langle U_{Lh}\rangle_{ei} \not\equiv 0$. As we will see below, the first possibility can be also used to explain the deviation of $R_K^*$, and $R_K$ from the SM prediction. These two conditions can be achieved by suitably choosing the mass matrices $m_D$ and $m_R$. Another possibility, allowing one to control the large effects in $\text{BR}(\mu \rightarrow e\gamma)$, which, anyway, we will not explore in this work, is to assume a $\zeta_\ell \not\equiv 0$ and tune it against the $\zeta_\ell^2$ term in Eq. (3.5) to reduce the $G_{H^\pm}$ form factor with respect to $G_{W^\pm}$.

With the analytic expressions of the Wilson coefficients, we may turn to finding realistic benchmark scenarios. We explore the parameter space spanned by the three parameters: $\zeta_\nu$, $m_{\nu_i}$ and $\langle U_{Lh}\rangle_{\mu,i}^2$ considering the impact of a single heavy neutrino, thus assuming, for the sake of simplicity, a hierarchy in the neutrino mass spectrum. The general case can be obtained straightforwardly and does not add much to the present discussion. In particular, one finds that if all the neutrino masses are almost degenerate, the combination $\sum_{i=1}^6 \langle U_{Lh}\rangle_{\mu,i}^2$ can be factored out from the Wilson coefficients and can be treated as an independent parameter leading to the same conclusions of the single neutrino case. This combination of squared mixing angles is bound from the non-unitarity test of the PMNS matrix to be less than $\sim 0.4 \times 10^{-3}$. Notice that we considered the scenario $\langle U_{Lh}\rangle_{i\mu} \not\equiv 0, \langle U_{Lh}\rangle_{ei} \simeq 0$ in line with the previous discussion on lepton flavour violating processes. The performed scan takes the ranges: $-80 < \zeta_\nu < 80$, $10^{-5} < \langle U_{Lh}\rangle_{\mu,i}^2 < 10^{-3}$ and 200 GeV $< m_{\nu_i} < 2000$ GeV. The other parameters are chosen as follows: $M_{H^\pm} = 550$ GeV, $\zeta_a = 1$ and $\zeta_d = \zeta_\ell \simeq 0$ which ensure that the flavour constraints discussed above, namely the ones which are not significantly affected by the presence of the heavy neutrinos, are all satisfied [26]. Other choices are obviously acceptable but not considered here.

The main result of the analysis is presented in Fig. 3. The numerical values of the flavour observables $R_K^*$ and $R_K$ have been obtained from the evaluation of the Wilson coefficients computed above and from the \texttt{flavio} package [27]. In the interesting region of the parameter space, the predictions for the two ratios in the A2HDM are the same and, for the sake of simplicity we will only discuss $R_K^*$. The latter is presented as a function of the parameter $\zeta_\nu Y_\nu \equiv \sqrt{2} \zeta_\nu \langle U_{Lh}\rangle_{i\mu} (m_{\nu_i}/v)$ that mostly controls the $C_{9,10}$
Figure 3. $R_{K^*}$ in the central bin and BR($B_s \to \mu^+\mu^-$) as a function of the combination of parameters that mostly controls the observable. The green bands are the 1σ and 2σ bands for the $R_{K^*}$ measurement. Blue (red) points correspond to the scenario in which the heavy neutrino has a non-negligible coupling only to the muon (electron) flavour eigenstates.

Wilson coefficients. In particular, the blue points correspond to the configuration in which the heavy neutrino couples to the SM muon sector and has negligible mixing with the first family. As anticipated above, this setup allows to reproduce the measured reduction in the $R_{K^*}$ ratio which is represented in the plot with a dashed horizontal line, together with 1σ and 2σ (green) bands. The red points, instead, are representative of the scenario in which the heavy neutrino has a mixing with the electrons. In this case the predicted $R_{K^*}$ is above one and contradicts the LHCb observations. The extent of the reduction is mainly controlled by the parameter $\zeta_\nu Y_\nu$. Interestingly, the same parameter defines the strength of the coupling among the charged Higgs, heavy neutrino and lepton and, as such, is subject to the perturbativity bound of [28]. The upper limit for the coupling is usually extracted from naive scaling arguments but also depends on the loop functions of the involved processes. In the plot, to facilitate the reading, we have shown two commonly adopted upper bounds.

Figure 4. (a) Correlation between $R_{K^*}$ in the central and low bins. (b) Correlation between $R_{K^*}$ in the central bin and the predicted BR($B_s \to \mu^+\mu^-$). Here only the 1σ bounds have been considered in both figures.
that can be used to estimate the region of tree-level perturbativity. With a $\zeta_\nu Y_\nu$ coupling below the most stringent perturbativity bound it is possible to explain the present $R_{K}$-measurement and the $\text{BR}(B_s \rightarrow \mu^+\mu^-)$ one within the 2$\sigma$ interval.

Fig. 4(a) displays the 1$\sigma$ bounds for both the $R_{K}$ low and central bins. One can see that parameter configurations can satisfy either the low or central bins separately, but not simultaneously both. In Fig. 4(b) we consider the effects on the $\text{BR}(B_s \rightarrow \mu^+\mu^-)$. One can see that for several points both predictions are simultaneously compatible with the experimental measurements.

6 Conclusions

The LHCb experiment at CERN has recently reported the existence of some anomalies in their data, with respect to the predictions of the SM. Specifically, the measured values of the observables $R_K = \text{BR}(B^+ \rightarrow K^+\mu^+\mu^-)/\text{BR}(B^+ \rightarrow K^+e^+e^-)$ and $R_{K^*} = \text{BR}(B^0 \rightarrow K^{*0}\mu^+\mu^-)/\text{BR}(B^0 \rightarrow K^{*0}e^+e^-)$ revealed a $\approx 2.5\sigma$ deviation when compared to the SM rates, which are essentially 1. In addition, the discrepancies occur in two di-lepton invariance mass bins. Therefore, these results must be taken as a serious hint of possible BSM physics.

Herein, we have considered the possibility of explaining such $R_K$ and $R_{K^*}$ anomalies in an A2HDM, wherein an alignment is present between the Yukawa couplings generated by the two Higgs doublet fields, combined with a low scale seesaw mechanism generating light neutrino masses and mixings in compliance with current experimental measurements. Such a scenario allows for significant non-universal leptonic contributions, through box diagrams mediated by $H^\pm$ and $\nu_R$ states, which in turn alter the yield of the partonic decay $b \rightarrow s\ell^+\ell^-$ entering the definition of both the $R_K$ and $R_{K^*}$ observables. In order to render our explanation phenomenologically viable, we have made sure to comply with both theoretical (chiefly, the unitarity bounds stemming from the neutrino mixing matrix) and experimental (the strongest being those due to lepton flavour violating processes and mesonic decay channels) constraints. In fact, the masses required for the charged Higgs and right-handed neutrino states entering the above transition are also well beyond their current direct limits. Furthermore, in order to make clear that our explanation of the $R_K$ and $R_{K^*}$ anomalies is not particularly ad hoc, we have left the actual low scale dynamics onsetting the seesaw mechanism undetermined, by illustrating that this could be realised through different scenarios, e.g., the so-called inverse and linear seesaw cases. Therefore, our setup captures a variety of light neutrino masses and mixings that can be tuned to further experimental observation in the neutrino sector while leaving predictions in the $B$ one unchanged. Finally, we have correlated our predictions for $R_K$ and $R_{K^*}$ in both di-lepton invariant mass bins to those for the highly constraining observable $\text{BR}(B_s \rightarrow \mu^+\mu^-)$ showing that a simultaneous solutions to both sets of measurements can be found in the envisioned A2HDM plus low scale seesaw scenario.
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