Solving the muon $(g-2)_\mu$ Anomaly in Two Higgs Doublet Models

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Summary. — Updating various theoretical and experimental constraints on the four different types of two-Higgs-doublet models (2HDMs), we find that only the "lepton-specific" (or "type X") 2HDM can explain the present muon (g-2) anomaly in the parameter region of large $\tan\beta$, a light CP-odd Higgs boson, and heavier CP-even and charged Higgs bosons which are almost degenerate. The severe constraints on the models come mainly from the consideration of vacuum stability and perturbativity, the electroweak precision data, B physics observables like $b \to s\gamma$ as well as the 125 GeV Higgs boson properties measured at the LHC.

PACS 12.60.Fr,13.40.Em 14.80.Bn, 14.80.Ec – .

1. – Outline

Since the first measurement of the muon anomalous magnetic moment $a_\mu = (g-2)_\mu/2$ by the E821 experiment at BNL in 2001 [1], much progress has been made in both experimental and theoretical sides to reduce the uncertainties by a factor of two or so establishing a consistent 3σ discrepancy

$$\Delta a_\mu \equiv a_\mu^{\rm EXP} - a_\mu^{\rm SM} = +262\,(85) \times 10^{-11} \tag{1}$$

which is in a good agreement with the different group's determinations. Since the 2001 announcement, there have been quite a few studies in the context of 2HDMs [2, 3, 4] restricted only to the type I and II models. However, the type X model [5] has some unique features in explaining the a_μ anomaly while evading all the experimental constraints.

Among many recent experimental results further confirming the Standard Model (SM) predictions, the discovery of the 125 GeV Brout-Egnlert-Higgs boson, which is very much SM-like, particularly motivates us to revisit the issue of the muon g$-$2 in favor of the type X 2HDM.

The key features in confronting 2HDMs with the muon g−2 anomaly can be summarized as follows [6, 7, 8, 9].

- The Barr-Zee two loop [10] can give a dominant (positive) contribution to the muon g−2 for a light CP-odd Higgs boson A and large $\tan\beta$ in the type II and X models.
- In the type II model, a light A has a large bottom Yukawa coupling for large $\tan\beta$, and thus is strongly constrained by the collider searches which have not been able to cover a small gap of 25 (40) GeV $< M_A <$ 70 GeV at the 2 (1) σ range of the muon (g-2) explanation [3].
- In the type II (and Y) model, the measured $\bar{B} \to X_s\gamma$ branching ratio pushes the charged Higgs boson $H^\pm$ high up to 480 (358) GeV at 95 (99) % C.L. [11], which requires a large separation between M_A and $M_{H^\pm}$ putting a strong limitation on the model due to the ρ parameter bound [4].
- Consideration of the electroweak precision data (EWPD) combined with the theoretical constraints from the vacuum stability and perturbativity requires the charged Higgs boson almost degenerate with the heavy Higgs boson H [12] (favoring $M_{H^\pm} > M_H$) and lighter than about 250 GeV in "the SM limit"; $\cos(\beta-\alpha) \to 0$. This singles out the type X model in favor of the muon g−2 [6].
- In the favored low m_A region, the 125 GeV Higgs decay $h \to AA$ has to suppressed kinematically or by suppressing the trilinear coupling λ_{hAA} which is generically order-one. This excludes the 1 σ range of the muon g−2 explanation in the SM limit [6].

However, the latest development [7, 8, 9] revealed more interesting possibilities in the "wrong-sign" domain (negative hbb or $h\tau\tau$ coupling) of 2HDMs [13].

- A cancellation in λ_{hAA} can be arranged to suppress arbitrarily the $h \to AA$ decay only in the wrong-sign limit with the heavy Higgs masses in the range of $M_{H^\pm} \sim M_H \approx 200 - 600$ GeV [7].
- The lepton universality affected by a large $H^+\tau\nu_\tau$ coupling turns out to severely constrain the large $\tan\beta$ and light $H^\pm$ region of the type X (and II) model and thus only a very low M_A and $\tan\beta$ region is allowed at 2 σ to explain the a_μ anomaly [8].

2. – Four types of 2HDMs

Non-observation of flavour changing neutral currents restricts 2HDMs to four different classes which differ by how the Higgs doublets couple to fermions [14]. They are organized by a discrete symmetry Z_2 under which different Higgs doublets and fermions carry different parities. These models are labeled as type I, II, "lepton-specific" (or X) and

TABLE I. – *The normalized Yukawa couplings for up- and down-type quarks and charged leptons.*

	y_u^A	y_d^A	y_l^A	y_u^H	y_d^H	y_l^H	y_u^h	y_d^h	y_l^h
Type I	$\cot\beta$	$-\cot\beta$	$-\cot\beta$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$
Type II	$\cot\beta$	$\tan\beta$	$\tan\beta$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\cos\beta}$	$\frac{\cos\alpha}{\cos\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$-\frac{\sin\alpha}{\cos\beta}$
Type X	$\cot\beta$	$-\cot\beta$	$\tan\beta$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\cos\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$
Type Y	$\cot\beta$	$\tan\beta$	$-\cot\beta$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\cos\beta}$	$\frac{\sin\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$\frac{\cos\alpha}{\sin\beta}$

"flipped" (or Y). Having two Higgs doublets $\Phi_{1,2}$, the most general Z_2 symmetric scalar potential takes the form:

$$
\begin{aligned}
V &= m_{11}^2|\Phi_1|^2 + m_{22}^2|\Phi_2|^2 - m_{12}^2(\Phi_1^\dagger\Phi_2 + \Phi_1\Phi_2^\dagger) \\
&+\frac{\lambda_1}{2}|\Phi_1|^4 + \frac{\lambda_2}{2}|\Phi_2|^4 + \lambda_3|\Phi_1|^2|\Phi_2|^2 + \lambda_4|\Phi_1^\dagger\Phi_2|^2 + \frac{\lambda_5}{2}\left[(\Phi_1^\dagger\Phi_2)^2 + (\Phi_1\Phi_2^\dagger)^2\right],
\end{aligned} \tag{2}
$$

where a (soft) Z_2 breaking term m_{12}^2 is introduced. Minimization of the scalar potential determines the vacuum expectation values $\langle\Phi_{1,2}^0\rangle \equiv v_{1,2}/\sqrt{2}$ around which the Higgs doublet fields are expanded as

$$
\Phi_{1,2} = \left[\eta_{1,2}^+, \frac{1}{\sqrt{2}}\left(v_{1,2} + \rho_{1,2} + i\eta_{1,2}^0\right)\right]. \tag{3}
$$

The model contains the five physical fields in mass eigenstates denoted by $H^\pm, A, H$ and h. Assuming negligible CP violation, $H^\pm$ and A are given by

$$
H^\pm, A = s_\beta\,\eta_1^{\pm,0} - c_\beta\,\eta_2^{\pm,0} \tag{4}
$$

where the angle β is determined from $t_\beta \equiv \tan\beta = v_2/v_1$, and their orthogonal combinations are the corresponding Goldstone modes $G^{\pm,0}$. The neutral CP-even Higgs bosons are diagonalized as

$$
h = c_\alpha\,\rho_1 - s_\alpha\,\rho_2, \quad H = s_\alpha\,\rho_1 + c_\alpha\,\rho_2 \tag{5}
$$

where $h\,(H)$ denotes the lighter (heavier) state.

The gauge couplings of h and H are given schematically by $\mathcal{L}_{\rm gauge} = g_V m_V\big(s_{\beta-\alpha}h + c_{\beta-\alpha}H\big)VV$ where $V = W^\pm$ or Z. When h is the 125 GeV Higgs boson, the SM limit corresponds to $s_{\beta-\alpha} \to 1$. Indeed, LHC finds, $c_{\beta-\alpha} \ll 1$ in all the 2HDMs confirming the SM-like property of the 125 GeV boson [15].

Normalizing the Yukawa couplings of the neutral bosons to a fermion f by m_f/v where $v = \sqrt{v_1^2+v_2^2} = 246$ GeV, we have the following Yukawa terms:

$$
\begin{aligned}
-\mathcal{L}_{\rm Yukawa}^{\rm 2HDMs} &= \sum_{f=u,d,l}\frac{m_f}{v}\left(y_f^h h\bar{f}f + y_f^H H\bar{f}f - iy_f^A A\bar{f}\gamma_5 f\right) \\
&+\left[\sqrt{2}V_{ud}H^+\bar{u}\left(\frac{m_u}{v}y_u^A P_L + \frac{m_d}{v}y_A^d P_R\right)d + \sqrt{2}\frac{m_l}{v}y_l^A H^+\bar{\nu}P_R l + h.c.\right]
\end{aligned} \tag{6}
$$

where the normalized Yukawa coupligs $y_f^{h,H,A}$ are summarized in Table I for each of these four types of 2HDMs.

Let us now recall that the tau Yukawa coupling $y_\tau \equiv y_l^h$ in Type X ($y_b \equiv y_d^h$ in Type II) can be expressed as

$$y_\tau = -\frac{s_\alpha}{c_\beta} = s_{\beta-\alpha} - t_\beta c_{\beta-\alpha} \tag{7}$$

which allows us to have the wrong-sign limit $y_\tau \sim -1$ compatible with the LHC data [13] if $c_{\beta-\alpha} \sim 2/t_\beta$ for large $\tan\beta$ favoured by the muon g−2. Later we will see that a cancellation in λ_{hAA} can be arranged only for $y_\tau^h < -1$ to suppress the $h \to AA$ decay.

3. – Electroweak constraints

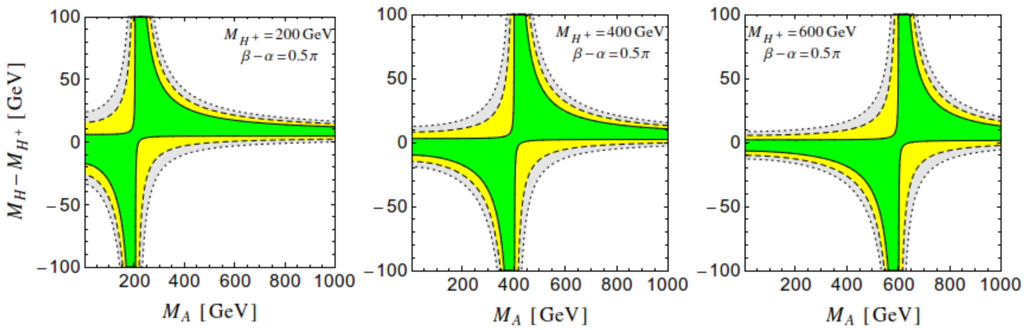

Fig. 1. – The parameter space allowed in the M_A vs. $\Delta M_H = M_H - M_{H^\pm}$ plane by EW precision constraints. The green, yellow, gray regions satisfy $\Delta\chi^2_{\rm EW}(M_A, \Delta M) < 2.3, 6.2, 11.8$, corresponding to 68.3, 95.4, and 99.7% confidence intervals, respectively.

Let us fist consider the constraints arising from EWPD on 2HDMs. In particular, we compare the theoretical 2HDMs predictions for M_W and $\sin^2\theta_{\rm eff}^{\rm lept}$ with their present experimental values via a combined χ^2 analysis. These quantities can be computed perturbatively by means of the following relations

$$M_W^2 = \frac{M_Z^2}{2}\left[1 + \sqrt{1 - \frac{4\pi\alpha_{\rm em}}{\sqrt{2}G_F M_Z^2}\frac{1}{1-\Delta r}}\right] \tag{8}$$

$$\sin^2\theta_{\rm eff}^{\rm lept} = k_l\left(M_Z^2\right)\sin^2\theta_W\,, \tag{9}$$

where $\sin^2\theta_W = 1 - M_W^2/M_Z^2$, and $k_l(q^2) = 1 + \Delta k_l(q^2)$ is the real part of the vertex form factor $Z \to l\bar{l}$ evaluated at $q^2 = M_Z^2$. We than use the following experimental values:

$$\begin{aligned} M_W^{\rm EXP} &= 80.385 \pm 0.015 \text{ GeV}, \\ \sin^2\theta_{\rm eff}^{\rm lept,EXP} &= 0.23153 \pm 0.00016. \end{aligned} \tag{10}$$

The results of our analysis are displayed in Fig. 1 confirming a custodial symmetry limit of our interest $M_A \ll M_H \sim M_{H^\pm}$ (or $M_H \ll M_A \sim M_{H^\pm}$) [12].

4. – Theoretical Constraints on the splitting M_A-M_{H^+}

Although any value of M_A is allowed by the EW precision tests in the limit of $M_H \sim M_{H^\pm}$, a large separation between $M_{H^\pm}$ and M_A is strongly constrained by theoretical requirements of vacuum stability, global minimum, and perturbativity:

$$\lambda_{1,2} > 0, \quad \lambda_3 > -\sqrt{\lambda_1\lambda_2}, \quad |\lambda_5| < \lambda_3 + \lambda_4 + \sqrt{\lambda_1\lambda_2}, \tag{11}$$

$$m_{12}^2(m_{11}^2 - m_{22}^2\sqrt{\lambda_1/\lambda_2})(\tan\beta - (\lambda_1/\lambda_2)^{1/4}) > 0, \tag{12}$$

$$|\lambda_i| \lesssim |\lambda_{\max}| = \sqrt{4\pi}, 2\pi, \text{ or } 4\pi. \tag{13}$$

Taking λ_1 as a free parameter, one can have the following expressions for the other couplings in the large t_β limit [9]:

$$\lambda_2 v^2 \approx s_{\beta-\alpha}^2 M_h^2 \tag{14}$$

$$\lambda_3 v^2 \approx 2M_{H^\pm}^2 - (s_{\beta-\alpha}^2 + s_{\beta-\alpha}y_\tau)M_H^2 + s_{\beta-\alpha}y_\tau M_h^2 \tag{15}$$

$$\lambda_4 v^2 \approx -2M_{H^\pm}^2 + s_{\beta-\alpha}^2 M_H^2 + M_A^2 \tag{16}$$

$$\lambda_5 v^2 \approx s_{\beta-\alpha}^2 M_H^2 - M_A^2 \tag{17}$$

where we have used the relation (7) neglecting the terms of $\mathcal{O}(1/t_\beta^2)$.

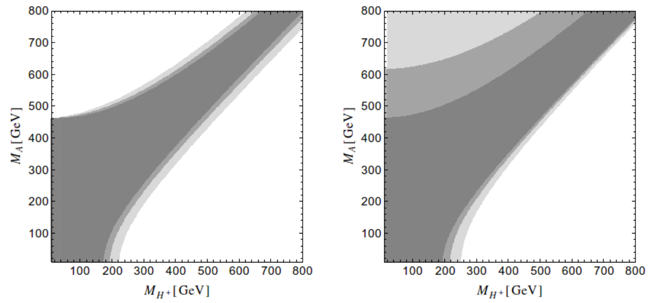

Fig. 2. – Theoretical constraints on the M_A-$M_{H^\pm}$ plane. The darker to lighter gray regions in the left panel correspond to the allowed regions for $\Delta M \equiv M_H - M_{H^\pm} = \{20, 0, -30\}$ GeV and $\lambda_{\max} = \sqrt{4\pi}$. The allowed regions in the right panel correspond to $\lambda_{\max} = \{\sqrt{4\pi}, 2\pi, 4\pi\}$ and vanishing ΔM.

Consideration of all the theoretical constraints mentioned above in the SM limit corresponding to $s_{\beta-\alpha} = y_\tau = 1$ gives us Fig. 2. One can see that for a light pseudoscalar with $M_A \lesssim 100$ GeV the charged Higgs boson mass gets an upper bound of $M_{H^\pm} \lesssim 250$ GeV.

5. – Constraints from the muon g−2

Considering all the updated SM calculations of the muon g−2, we obtain

$$a_\mu^{\rm SM} = 116591829\,(57) \times 10^{-11} \tag{18}$$

comparing it with the experimental value $a_\mu^{\rm EXP} = 116592091\,(63) \times 10^{-11}$, one finds a deviation at 3.1σ: $\Delta a_\mu \equiv a_\mu^{\rm EXP} - a_\mu^{\rm SM} = +262\,(85) \times 10^{-11}$. In the 2HDM, the one-loop contributions to a_μ of the neutral and charged Higgs bosons are

$$\delta a_\mu^{\rm 2HDM}(1{\rm loop}) = \frac{G_F\, m_\mu^2}{4\pi^2\sqrt{2}} \sum_j \left(y_\mu^j\right)^2 r_\mu^j\, f_j(r_\mu^j), \tag{19}$$

where $j = \{h, H, A, H^\pm\}$, $r_\mu^j = m_\mu^2/M_j^2$, and

$$f_{h,H}(r) = \int_0^1 dx\, \frac{x^2(2-x)}{1-x+rx^2}, \tag{20}$$

$$f_A(r) = \int_0^1 dx\, \frac{-x^3}{1-x+rx^2}, \tag{21}$$

$$f_{H^\pm}(r) = \int_0^1 dx\, \frac{-x(1-x)}{1-(1-x)r}. \tag{22}$$

These formula show that the one-loop contributions to a_μ are positive for the neutral scalars h and H, and negative for the pseudo-scalar and charged Higgs bosons A and $H^\pm$ (for $M_{H^\pm} > m_\mu$). In the limit $r \ll 1$,

$$f_{h,H}(r) = -\ln r - 7/6 + O(r), \tag{23}$$

$$f_A(r) = +\ln r + 11/6 + O(r), \tag{24}$$

$$f_{H^\pm}(r) = -1/6 + O(r), \tag{25}$$

showing that in this limit $f_{H^\pm}(r)$ is suppressed with respect to $f_{h,H,A}(r)$. Now the two-loop Barr-Zee type diagrams with effective $h\gamma\gamma$, $H\gamma\gamma$ or $A\gamma\gamma$ vertices generated by the exchange of heavy fermions gives

$$\delta a_\mu^{\rm 2HDM}(2{\rm loop}-{\rm BZ}) = \frac{G_F\, m_\mu^2}{4\pi^2\sqrt{2}}\, \frac{\alpha_{\rm em}}{\pi} \sum_{i,f} N_f^c\, Q_f^2\, y_\mu^i\, y_f^i\, r_f^i\, g_i(r_f^i), \tag{26}$$

where $i = \{h, H, A\}$, $r_f^i = m_f^2/M_i^2$, and m_f, Q_f and N_f^c are the mass, electric charge and number of color degrees of freedom of the fermion f in the loop. The functions $g_i(r)$ are

$$g_i(r) = \int_0^1 dx\, \frac{\mathcal{N}_i(x)}{x(1-x)-r} \ln \frac{x(1-x)}{r}, \tag{27}$$

where $\mathcal{N}_{h,H}(x) = 2x(1-x) - 1$ and $\mathcal{N}_A(x) = 1$.

Note the enhancement factor m_f^2/m_μ^2 of the two-loop formula in Eq. (26) relative to the one-loop contribution in Eq. (19), which can overcome the additional loop suppression factor α/π, and makes the two-loop contributions may become larger than the one-loop ones. Moreover, the signs of the two-loop functions $g_{h,H}$ (negative) and g_A (positive) for the CP-even and CP-odd contributions are opposite to those of the functions $f_{h,H}$ (positive) and f_A (negative) at one-loop. As a result, for small M_A and large $\tan\beta$ in Type II and X, the positive two-loop pseudoscalar contribution can generate a dominant contribution which can account for the observed Δa_μ discrepancy. The additional 2HDM contribution $\delta a_\mu^{\rm 2HDM} = \delta a_\mu^{\rm 2HDM}(1{\rm loop}) + \delta a_\mu^{\rm 2HDM}(2{\rm loop}-{\rm BZ})$ obtained adding Eqs. (19) and (26) (without the h contributions) is compared with Δa_μ in Fig. 3.

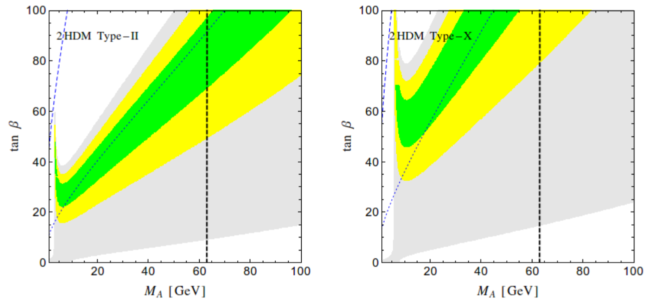

Fig. 3. – The 1σ, 2σ and 3σ regions allowed by Δa_μ in the M_A-$\tan\beta$ plane taking the limit of $\beta-\alpha=\pi/2$ and $M_{h(H)} = 126$ (200) GeV in type II (left panel) and type X (right panel) 2HDMs. The regions below the dashed (dotted) lines are allowed at 3σ (1.4σ) by Δa_e. The vertical dashed line corresponds to $M_A = M_h/2$.

Finally, let us remark that the hAA coupling is generically order one and thus can leads to a sizable non-standard decay of $h\to AA$ which should be suppressed kinematically or by making $|\lambda_{hAA}/v| \ll 1$ to meet the LHC results [9, 7, 8]. Using Eq. (14), one gets the hAA coupling, $\lambda_{hAA}/v \approx s_{\beta-\alpha}[\lambda_3+\lambda_4-\lambda_5]$, and thus

$$\lambda_{hAA} v/s_{\beta-\alpha} \approx -(1+s_{\beta-\alpha}y_\tau)M_H^2 + s_{\beta-\alpha}y_\tau M_h^2 + 2M_A^2 \tag{28}$$

where we have put $s^2_{\beta-\alpha}=1$ [9]. It shows that, in the SM limit of $s_{\beta-\alpha}y_\tau \to 1$, the condition $\lambda_{hAA}\approx 0$ requires $M_H \sim M_h$ which is disfavoured, and thus one needs to have $M_A > M_h/2$. On the other hand, one can arrange a cancellation for $\lambda_{hAA}\approx 0$ in the wrong-sign domain $s_{\beta-\alpha}y_\tau < 0$ if the tau Yukawa coupling satisfies

$$y_\tau s_{\beta-\alpha} \approx -\frac{M_H^2 - 2M_A^2}{M_H^2 - M_h^2}. \tag{29}$$

6. – Summary

The type X 2HDM provides a unique opportunity to explain the current $\sim 3\sigma$ deviation in the muon g$-$2 while satisfying all the theoretical requirements and the experimental constraints. The parameter space favourable for the muon g-2 at 2σ is quite limited in the SM limit: $\tan\beta \gtrsim 30$ and $M_A \ll M_H \sim M_{H^\pm} \lesssim 250$ GeV. However, consideration of the $h \to AA$ decay and lepton universality [8] rules out this region. On the other hand, in the wrong-sign limit of $y_\tau \sim -1$, a cancellation for $\lambda_{hAA} \approx 0$ can be arranged for M_H up to about 600 GeV [7, 9] opening up more parameter space.

Such a light CP-odd boson A and the extra Heavy bosons can be searched for at the next run of the LHC mainly through $pp \to AH, AH^\pm$ followed by the decays $H^\pm \to \tau^\pm \nu$ and $A, H \to \tau^+\tau^-$ [8, 9].

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