Cornering Spontaneous *CP* Violation with Charged-Higgs-Boson Searches

Ulrich Nierste,* Mustafa Tabet,† and Robert Ziegler‡
Institut für Theoretische Teilchenphysik (TTP), Karlsruher Institut für Technologie (KIT), 76131 Karlsruhe, Germany

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Decades of precision measurements have firmly established the Kobayashi-Maskawa phase as the dominant source of the charge-parity (CP) violation observed in weak quark decays. However, it is still unclear whether CP violation is explicitly encoded in complex Yukawa matrices or instead stems from spontaneous symmetry breaking with underlying CP-conserving Yukawa and Higgs sectors. Here we study the latter possibility for the case of a generic two-Higgs-doublet model. We find that theoretical constraints limit the ratio t_β of the vacuum expectation values (vevs) to the range $0.22 \leq t_\beta \leq 4.5$ and imply the upper bounds $M_{H^\pm} \leq 435$ GeV, $M_{H_2^0} \leq 485$ GeV and $M_{H_3^0} \leq 545$ GeV for the charged and extra neutral Higgs masses. We derive lower bounds on charged-Higgs couplings to bottom quarks which provide a strong motivation to study the nonstandard production and decay signatures $pp \to qbH^\pm(\to q'b)$ with *all* flavors $q, q' = u, c, t$ in the search for the charged Higgs boson. We further present a few benchmark scenarios with interesting discovery potential in collider analyses.

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Introduction.—In 1964 the observation of the decay $K_L \to \pi\pi$ has established the violation of charge-parity (CP) symmetry [1]. Owing to the CPT theorem [2] this discovery implies that also time-reversal symmetry (T) is broken and nature has a microscopic arrow of time. In 1973 two landmark papers proposed possible mechanisms of CP violation (CPV) involving new particles: M. Kobayashi and T. Maskawa (KM) pointed out that explicit CPV can occur if the Standard Model (SM) is amended by a third quark generation [3], while T. D. Lee showed that spontaneous CPV can be realized in the presence of a second Higgs doublet [4].

The subsequent success of the KM mechanism, however, did not rule out the possibility of spontaneous CP violation: The complex phase of the Cabibbo-Kobayashi-Maskawa (CKM) matrix stems from the diagonalization of complex quark mass matrices, and these matrices may still arise as linear combinations of real Yukawa matrices multiplied by complex vevs.

Almost half a century later, the issue of explicit vs spontaneous CPV still remains unresolved! The main purpose of this paper is to tackle this question systematically and discuss how to either discover spontaneous CPV or to entirely rule out this possibility using future data from precision observables and colliders. The latter is possible, because spontaneous-CPV scenarios have no decoupling limit and feature a pattern of flavor violation that cannot be aligned to the SM.

The main obstacle to this endeavor is the considerable size of the parameter space of SCPV models. Indeed previous works have so far considered only special cases of two-Higgs-doublet model (2HDM) (see, e.g., Refs. [5,6]). Our Letter targets generic features of SCPV and only makes two simplifying assumptions, which are justified by shortcutting to that region of the parameter space that is least constrained by experiment. First, we identify the lightest neutral Higgs boson with the 125 GeV SM-like Higgs particle. Second, we do not permit Yukawa terms leading to FCNC Higgs couplings among down-type quarks, which are severely constrained by precision flavor data. These data constrain the mentioned couplings so tightly that relaxing our second assumption will not change our conclusions. We find a remarkable sum rule for charged-Higgs couplings to b quarks, which implies that at least one of the couplings to tb, cb, or ub is sizable. Given the upper limit on the charged Higgs mass and the constraints from precision observables, these results reveal that charged Higgs searches in non-standard channels have the potential to either support or falsify SCPV as the primary origin of the KM phase.

General features.—Higgs sector: The most general potential with two SU(2) Higgs doublet fields $\phi_i = (\phi_i^0, \phi_i^+)^T$, $i = 1, 2$, reads [4]

$$V = m_{11}^2\phi_1^\dagger\phi_1 + m_{22}^2\phi_2^\dagger\phi_2 - (m_{12}^2\phi_1^\dagger\phi_2 + \mathrm{H.c.}) + \frac{\lambda_1}{2}(\phi_1^\dagger\phi_1)^2 + \frac{\lambda_2}{2}(\phi_2^\dagger\phi_2)^2 + \lambda_3(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) + \lambda_4(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) + \left[\phi_1^\dagger\phi_2\left(\frac{\lambda_5}{2}\phi_1^\dagger\phi_2 + \lambda_6(\phi_1^\dagger\phi_1) + \lambda_7(\phi_2^\dagger\phi_2)\right) + \mathrm{H.c.}\right], \tag{1}$$

Adopting canonical CP transformation rules, $CP\phi_i(x^\mu) = \phi_i^*(x_\mu)$, CP conservation means that all parameters in Eq. (1) are real. For appropriate choices of these parameters V is bounded from below and has a local minimum for the complex vevs:

$$\langle\phi_1\rangle = \begin{pmatrix} 0 \\ vc_\beta \end{pmatrix}, \qquad \langle\phi_2\rangle = \begin{pmatrix} 0 \\ vs_\beta e^{i\xi} \end{pmatrix}, \tag{2}$$

with $c_\beta \equiv \cos\beta > 0$, $s_\beta \equiv \sin\beta > 0$, $v = 174$ GeV, and the CP phase ξ. This minimum is in fact always the global one, which immediately follows from the results in Ref. [7]. As an important observation, the three minimization equations with respect to $\mathrm{Re}\phi_1^0$, $\mathrm{Re}\phi_2^0$, and $\mathrm{Im}\phi_2^0$ allow to trade all three massive parameters m_1^2, m_2^2, m_{12}^2 in Eq. (1) for the three vev parameters v, $t_\beta \equiv \tan\beta$ and ξ. Therefore all elements of the Higgs mass matrices are of the order of the electroweak scale v, with dimensionless coefficients composed of λ_{1-7}, $c_\xi \equiv \cos\xi$, $s_\xi \equiv \sin\xi$ and t_β. Since perturbativity does not permit arbitrarily large couplings, the masses of all Higgs bosons are bounded from above. This absence of a decoupling limit in the 2HDM with SCPV has been observed already in Refs. [6,8], for other examples see Refs. [9–11].

The Higgs spectrum consists of a charged Higgs with mass

$$m_{H^\pm} = v\sqrt{\lambda_5 - \lambda_4}, \tag{3}$$

and three neutral Higgs states H_A with masses that fulfill the sum rule

$$\frac{1}{2}\sum_{A=1}^{3}\frac{m_{H_A}^2}{v^2} = s_{2\beta}c_\xi(\lambda_6 + \lambda_7) + \lambda_2 s_\beta^2 + \lambda_1 c_\beta^2 + \lambda_5. \tag{4}$$

Requiring NLO perturbative unitarity [12] allows to derive upper bounds for the physical Higgs masses, which can be further tightened by identifying the lightest Higgs with the SM Higgs [13]. Using the results in Refs. [14–16], we find

$$m_{H^\pm} \lesssim 435 \text{ GeV}, \tag{5}$$

while neutral Higgs masses must satisfy

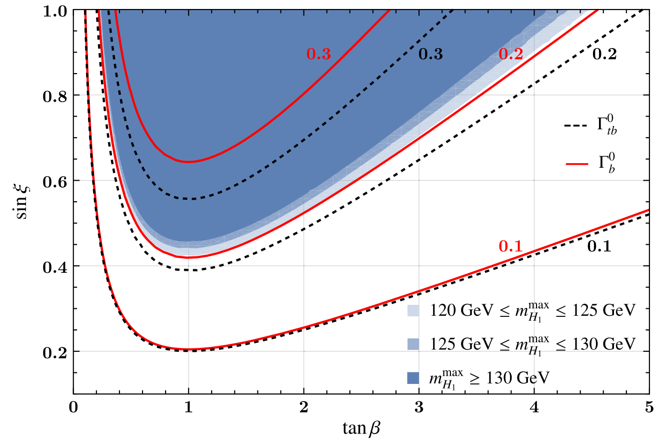

FIG. 1. Contours of Γ_b^0 (red, solid) and Γ_{tb}^0 (black, dashed) in the $t_\beta - s_\xi$ plane. We also indicate regions with different values of $m_{H_1}^{\max}$, which is the maximal mass for the tree-level value of the lightest Higgs H_1 allowed by NLO perturbativity.

$$m_{H_2} \lesssim 485 \text{ GeV}, \qquad m_{H_3} \lesssim 545 \text{ GeV}, \tag{6}$$

with the sum of all three neutral Higgs masses bounded by 1.1 TeV. Our bounds are tighter than those in Ref. [6], because we include the NLO corrections of Ref. [12]. Moreover, since the determinant of the neutral Higgs mass matrix is proportional to $s_\xi^2 s_{2\beta}^2$, requiring that all states are heavier than 125 GeV gives lower bounds on s_ξ and a range for t_β. Using again NLO perturbative unitarity, we find (see Fig. 1)

$$0.22 \lesssim t_\beta \lesssim 4.5, \qquad |s_\xi| \gtrsim 0.42. \tag{7}$$

The neutral Higgs mass basis is obtained by diagonalizing $O^T M_H^2 O = M_{H,\mathrm{diag}}^2$ with the orthogonal matrix

$$O \equiv R_{12}(\theta_{12} - \beta)R_{13}(\theta_{13})R_{23}(\theta_{23}), \tag{8}$$

where $R_{ij}(\theta)$ are rotation matrices in the $i - j$ plane by an angle θ_{ij}. Since the Higgs mass matrices only depend on λ_{1-7} (besides s_ξ, t_β, and v), we can trade the seven λ_i parameters for the four Higgs masses $m_{H^\pm}$, m_{H_i} and the three mixing angles $s_{ij} \equiv \sin\theta_{ij}$. These mixing angles appear in all couplings of the neutral Higgs mass eigenstates. The couplings to massive gauge bosons $g_{H_A VV}$ are given in terms of the corresponding SM Higgs couplings g_{hVV} by

$$g_{H_A VV} = (c_\beta O_{1A} + s_\beta O_{2A})g_{hVV}. \tag{9}$$

Particularly simple are the couplings of the lightest neutral Higgs $g_{H_1VV}/g_{hVV} = c_{12}c_{13}$. Since throughout this Letter we will assume that H_1 is the observed SM-like Higgs state with a mass of 125 GeV, its couplings need to be sufficiently close to the couplings of the SM Higgs, i.e., $s_{12}, s_{13} \ll 1$.

Yukawa sector: The quark Yukawa Lagrangian is given by

$$\mathcal{L}_{\rm yuk} = -\bar{Q}_L(Y_{u1}\tilde{\phi}_1 + Y_{u2}\tilde{\phi}_2)u_R - \bar{Q}_L(Y_{d1}\phi_1 + Y_{d2}\phi_2)d_R + \text{H.c.}, \tag{10}$$

with the Yukawa matrices Y_{qi} and $\tilde{\phi}_i = \epsilon_{ij}\phi_j^*$, $\epsilon_{12} = 1$. Since $\mathcal{L}_{\rm yuk}$ conserves CP, we can choose $Y_{d1,d2}$ real. This implies that fermion mass matrices, given by

$$\frac{M_u}{v} = Y_{u1}c_\beta + Y_{u2}e^{-i\xi}s_\beta, \qquad \frac{M_d}{v} = Y_{d1}c_\beta + Y_{d2}e^{i\xi}s_\beta, \tag{11}$$

can induce the KM phase only if ξ is physical, i.e., cannot be rotated away by field redefinitions. This implies flavor misalignment, defined through $Y_{q1}Y_{q2}^T - Y_{q2}Y_{q1}^T \neq 0$, which necessarily induces FCNC couplings of neutral Higgs bosons. Since Eq. (6) forbids arbitrarily heavy neutral Higgs bosons, one cannot suppress all Higgs-mediated FCNC processes simultaneously to arbitrarily small values. As constraints on FCNC Higgs couplings to down-type quarks are particularly strong, in the following we set $Y_{d2} \approx 0$, thus relegating all FCNC couplings to the up sector. We stress that this choice is dictated solely by phenomenological constraints, and note that it is radiatively stable, since loops corrections $\delta Y_{d2} \propto Y_{u1}Y_{u2}^T Y_{d1}/(16\pi^2)$ are numerically irrelevant.

Without loss of generality, we can work in a flavor basis where Y_{d1} is diagonal and $M_u = V^\dagger m_u^{\rm diag}V_R^\dagger$, where V is the CKM matrix and V_R a free unitary matrix. The Higgs couplings to fermions in the mass basis are then given by

$$\begin{aligned}\mathcal{L}_H = &-\bar{u}_{L,i}\frac{H_A^0}{\sqrt{2}}[\delta_{ij}\alpha_A^u + t_\beta\tilde{\epsilon}_{ij}^u\beta_A^u]\frac{m_{u_j}}{vs_\beta}u_{R,j}\\ &-\bar{d}_{L,i}\frac{\alpha_A^d H_A^0}{\sqrt{2}}\frac{m_{d_i}}{vc_\beta}d_{R,i}\\ &+\bar{d}_{L,i}H^-V_{ki}^*\left[\delta_{kj}c_\beta - \frac{\tilde{\epsilon}_{kj}^u}{c_\beta}\right]\frac{m_{u_j}}{vs_\beta}u_{R,j}\\ &+\bar{u}_{L,i}H^+V_{ij}s_\beta\frac{m_{d_j}}{vc_\beta}d_{R,j} + \text{H.c.},\end{aligned} \tag{12}$$

with

$$\alpha_A^u = O_{2A} - ic_\beta O_{3A}, \qquad \beta_A^u = O_{1A} - \frac{O_{2A}}{t_\beta} + i\frac{O_{3A}}{s_\beta},$$
$$\alpha_A^d = O_{1A} - is_\beta O_{3A}, \qquad \tilde{\epsilon}_{ij}^u = (VY_{u1}V_R)_{ij}\frac{vc_\beta}{m_{u_j}}. \tag{13}$$

Using Eq. (11), we can write Y_{u1} as

$$Y_{u1} = \frac{1}{vc_\beta}\left(\text{Re} + \frac{c_\xi}{s_\xi}\text{Im}\right)[V^\dagger m_u^{\rm diag}V_R^\dagger], \tag{14}$$

which entails an expression for the couplings $\tilde{\epsilon}_{ij}^u$:

$$\tilde{\epsilon}_{jk}^u = \frac{t_\xi - i}{2t_\xi}\delta_{jk} + \frac{t_\xi + i}{2t_\xi}(VV^T m_u V_R^T V_R m_u^{-1})_{jk}. \tag{15}$$

Note that if we use the residual re-phasing freedom to bring the CKM matrix to the usual Particle Data Group (PDG) convention $V_{\rm PDG}$, we have $V \to V_{\rm PDG}$ in Eq. (12), but $V \to V_{\rm PDG}P$ in Eq. (15) with a free (diagonal) phase matrix P. The Higgs couplings only depend on the combination $V_RV_R^T$, which in this phase convention is a generic symmetric unitary matrix with three physical phases. Apart from the angles and phases in V_R, all quark flavor violation in the Higgs sector is entirely determined by up-quark masses and CKM elements.

Taking the lepton Yukawa sector analogous to the down-quark sector, with only one Higgs doublet coupling to right-handed charged leptons, one obtains a SM-like phenomenology of charged-lepton decays. The $H^+\bar{\nu}_{\tau L}\tau_R$ coupling can neither vanish nor be much larger than m_τ/v, implied by the t_β range in Eq. (7).

Charged Higgs couplings: Since neutral Higgs couplings are more sensitive to the free parameters in V_R, we instead focus on the fermion couplings of the charged Higgs. Indeed, the peculiar structure of the Yukawa sector guarantees that at least one coupling of the charged Higgs to bottom quarks, $H^+\bar{u}_{iR}\Gamma_{u_ib}^{RL}b_L$, is sizable. Using Eq. (15) and unitarity of V_R, one can show that these couplings satisfy the remarkable relation

$$\sum_{i=u,c,t}|\Gamma_{ib}^{RL}|^2 = \frac{m_t^2}{v^2} + \frac{2m_t}{vs_{2\beta}}\left(c_{2\beta}^2\text{Re}\Gamma_{tb}^{RL} - \frac{\text{Im}\Gamma_{tb}^{RL}}{t_\xi}\right) + \mathcal{O}\left(|V_{cb}|\frac{m_c}{m_t}\right), \tag{16}$$

which follows solely from SCPV and the assumption that Y_{d2} is approximately diagonal in the same basis as Y_{d1}. This relation implies that the largest coupling $\Gamma_b^{\rm max} \equiv \max\{|\Gamma_{ub}^{RL}|, |\Gamma_{cb}^{RL}|, |\Gamma_{tb}^{RL}|\}$ is bounded from below

$$\Gamma_b^{\rm max} \geq \Gamma_b^0 \equiv \frac{A}{2n}(\sqrt{1+n\kappa} - 1), \tag{17}$$

where $n = 3$ and the rhs is only a function of β and ξ

$$A = \frac{2m_t}{vs_{2\beta}}\sqrt{c_{2\beta}^2 + 1/t_\xi^2}, \qquad \kappa = \frac{s_{2\beta}^2 t_\xi^2}{1 + c_{2\beta}^2 t_\xi^2}. \tag{18}$$

Minimizing Γ_b^0 over β and ξ as allowed by NLO perturbativity and $m_{H_1} \geq 120$ GeV, one numerically finds

$$\max\{|\Gamma_{ub}^{RL}|, |\Gamma_{cb}^{RL}|, |\Gamma_{tb}^{RL}|\} \geq \Gamma_b^0 \geq 0.20. \quad (19)$$

We show contours of Γ_b^0 in the $t_\beta - s_\xi$ plane in Fig. 1. As one can see from this plot, our lower bound on Γ_b^0 in Eq. (19) is rather conservative.

Note that $\Gamma_b^{\max}$ reaches its minimum Γ_b^0 for equal couplings $|\Gamma_{ib}^{LR}|$, i.e., if $\Gamma_b^{\max} = |\Gamma_{ib}^{LR}| = \Gamma_b^0$ for $i = u$, c, t. It is instructive to consider two other special cases: if $\Gamma_{ub}^{LR} = \Gamma_{cb}^{LR} = 0$, then $\Gamma_b^{\max}$ coincides with $|\Gamma_{tb}^{LR}|$ and has a minimal value Γ_{tb}^0 that is given by the rhs of Eq. (17), but with $n = 1$. Note that typically $\kappa \ll 1$, which implies that Γ_{tb}^0 is only slightly larger than Γ_b^0. The contours of Γ_{tb}^0 in the $t_\beta - s_\xi$ plane are also shown in Fig. 1, and indeed coincide with those of Γ_b^0 when Γ_b^0 and therefore κ are small. If instead $|\Gamma_{tb}^{LR}| = 0$, the couplings to light generations become large, since in this case they satisfy the sum rule $|\Gamma_{ub}^{LR}|^2 + |\Gamma_{cb}^{LR}|^2 = m_t^2/v^2$, which directly follows from Eq. (16).

The lower bound on charged Higgs couplings to b quarks in Eq. (19), together with the upper bound on the charged Higgs mass in Eq. (5) render our class of models predictive despite the considerable number of free parameters and Eq. (19) entails a “no-lose” theorem for charged-Higgs discovery.

Phenomenology.—The phenomenology depends on 17 free parameters: the heavy Higgs masses $m_{H^\pm}, m_{H_2}, m_{H_3}$, the vacuum angles β and ξ, the mixing angles s_{12}, s_{13}, s_{23} and three angles plus six phases that determine $\tilde{\epsilon}_u$ in Eq. (15). Although huge, this parameter space is compact because of the absence of new mass scales and perturbative unitarity, cf. Eqs. (5)–(7), which allows us to confirm or rule out the model in the near future. In the following we discuss indirect searches via precision measurements and direct searches for the new additional Higgs states.

The SM-like measurements of Higgs coupling strengths [17,18] imply small values of s_{13} and s_{12}, i.e., a Higgs sector close to the alignment limit. Also constraints from precision observables like neutral meson mixing [19–21], $B \to X_s\gamma$ [22], and electric dipole moments (EDMs) [23,24] have considerable impact on the parameter space, but do not exclude all of it. Indeed in certain, nontrivial parameter ranges all heavy Higgs couplings to fermions can be simultaneously suppressed to a level that all observables are SM-like. Still many precision observables can be close to their current experimental limits, for example electron and neutron EDMs can be as large as $|d_e| = 10^{-29}$ e cm and $|d_n| = 3 \times 10^{-26}$ e cm, respectively. These regions will be explored by several near-future experiments, like nEDM [25,26], n2EDM [27], and the eEDM experiment by the ACME collaboration [24]. Thus precision measurements will continue to probe the parameter space from below, pushing up the limits on heavy Higgs masses towards the unitarity limits in Eqs. (5) and (6).

Also present experimental data from direct Higgs searches constrain significant portions of the parameter space, but do not allow to exclude the entire scenario. Actually it is quite easy to evade standard searches while predicting sizable production cross sections for signatures that have not been looked for so far, in particular those that result from the dominance of flavor-violating Higgs couplings. Indeed it follows from the bound in Eq. (19) that the charged Higgs is guaranteed to have sizable couplings to bottom and up, charm, or top quarks. As charged Higgs couplings to gauge bosons are suppressed in the alignment limit, while couplings to leptons are bounded by the smallness of t_β, the quark couplings typically dominate both production and decay. In the following we briefly discuss the resulting charged Higgs phenomenology at the LHC, using the benchmark points in Table I as an illustration. A much more detailed analysis of the collider phenomenology will be presented elsewhere. Because of the upper limit in Eq. (5), we are only interested in the light mass range below 440 GeV, which is typically more difficult to probe at hadron colliders due to large SM backgrounds.

The case of tb associated production and decay to tb, $pp \to tbH^\pm(\to tb)$, belongs to the standard charged Higgs searches by CMS and ATLAS, cf. Refs. [30,31] and [32,33] for 8 TeV and 13 TeV data, respectively. These searches exclude signal strengths of $\mathcal{O}(1\text{ pb})$ in the relevant mass range. An exemplary benchmark point close to exclusion is provided by BP1 in Table I. Charged Higgs couplings to tb can be suppressed if couplings to cb or/and ub are enhanced, which corresponds to fairly unexplored signatures. The phenomenology of the case of cb dominance is extensively discussed in Ref. [34] (see also Ref. [35]). Apart from larger production cross sections and possible charm tagging in charged Higgs decays, the case of ub is quite similar to the one of cb, so we will focus on these cases in the remaining discussion, largely following Ref. [34]. A benchmark point with large ub coupling is provided in Table I by BP2.

Starting with $pp \to cbH^\pm(\to cb)$, this process can be probed at the LHC by inclusive searches for low-mass dijet resonances like Ref. [36], which however are not yet sensitive to charged Higgs masses below 450 GeV. Our scenario hopefully motivates further efforts to optimize future searches for resonances in multijet final states going towards lower masses. For example, we find benchmarks with (inclusive) production cross sections of $pp \to b(c)H^\pm$ as large as $\mathcal{O}(\text{nb})$, which are not excluded by present data, see BP3.

The next possibility is $pp \to cbH^\pm(\to tb)$, which is represented by BP4. Despite the large production cross sections of $\mathcal{O}(10\text{ pb})$ (for the case of untagged charm jets), experimental searches are hampered by the fact that the jets from the associated b and c quarks typically fall outside the trigger range for rapidity and transverse momentum (see however Ref. [37] for a recent study of the discovery potential using associated b jets). Thus only searches for tb

TABLE I. Table with benchmark points compatible with theoretical constraints and current experimental data. The total production cross section at 14 TeV $\sigma_{\rm prod}^{14\,\rm TeV}$ in the respective channel has been obtained with MadGraph5_aMC@NLO [28] using the 2HDM_NLO model file [29] with the default run cards. The cross sections for $pp \to qbH^{\pm}$ denote the sum of the inclusive cross sections for a final state containing an untagged $q = u$, c jet.

	BP1	BP2	BP3	BP4	BP5	BP6
m_{H_2} (GeV)	170	335	215	320	160	135
m_{H_3} (GeV)	245	355	245	335	190	335
$m_{H^{\pm}}$ (GeV)	180	375	170	325	165	350
t_{β}	0.38	0.84	0.43	0.67	0.36	0.37
ξ	1.95	1.96	1.59	1.59	2.04	1.32
$\lvert s_{13}\rvert \times 10^2$	4.4	1.3	0.12	0.081	7.5	1.2
$\lvert s_{12}\rvert \times 10^2$	2.3	1.6	0.32	0.095	2.8	3.4
$\lvert s_{23}\rvert$	0.23	1.00	0.18	0.99	0.42	0.21
$\lvert\Gamma_{tb}^{RLH^{\pm}}\rvert$	0.38	0.58	0.16	1.17	0.35	0.32
$\lvert\Gamma_{cb}^{RLH^{\pm}}\rvert$	$< 10^{-3}$	0.76	0.76	0.70	$< 10^{-2}$	0.37
$\lvert\Gamma_{ub}^{RLH^{\pm}}\rvert$	$< 10^{-6}$	0.45	$< 10^{-4}$	$< 10^{-5}$	$< 10^{-4}$	$< 10^{-7}$
$\sigma_{\rm prod}^{14\,\rm TeV}(pp \to qbH^{\pm})$ (pb)	≈ 0	190	520	36	≈ 0	7.9
$\sigma_{\rm prod}^{14\,\rm TeV}(pp \to tbH^{\pm})$ (pb)	0.62	0.28	0.13	1.6	0.71	0.10
$\sigma_{\rm Drell\text{-}Yan}^{14\,\rm TeV}(pp \to H^{\pm}H^{\mp})$ (fb)	35	2.0	44	3.7	49	3.5
Main decay channel of $H^{\pm}$	tb 99%	cb 58%	cb 100%	tb 59%	cb 89%	$WH_2(\to c\bar{c}, b\bar{b})$ 62%

resonances can be used, which at present focus on the heavy mass range above 1 TeV (see, e.g., Ref. [38]), and it is unclear whether further data and optimization will probe masses as low as 300 GeV.

Occasionally $pp \to tbH^{\pm}(\to cb)$ can be the main production and decay for charged Higgs masses that are close to the top threshold, see BP5. The signature is the same as in LHC searches for $tt \to WbH^{\pm}(\to cb)b$, which so far have been analyzed only for charged Higgs masses much below the top threshold, see, e.g., Ref. [39]. Thus our scenario motivates searches also for masses as large as 170 GeV, together with the models considered in Refs. [34,35].

Other possible signatures like charged Higgs decays into WH_2 depend on the details of heavy neutral Higgs phenomenology, which is more model dependent. Nevertheless we provide one benchmark point BP6 with dominant $H^{\pm} \to W^{\pm}H_2$ decay, where H_2 further decays to $\bar{c}c$ or $\bar{b}b$. Finally we note that also charged Higgs pair production via Drell-Yan provides a model-independent production channel that varies between 2 and 50 fb for the benchmark points in Table I.

Summary and conclusions.—We have discussed the generic framework of spontaneous CP violation in the 2HDM, where the KM phase arises solely from the Higgs potential. This scenario has the remarkable feature that all mass scales are set by the electroweak scale up to quartic couplings, so that perturbative unitarity implies model-independent upper bounds on all heavy Higgs states, cf. Eqs. (5) and (6). Moreover, the new scalar states must necessarily have a particular, nonstandard pattern of flavor violation in order to induce a nonvanishing KM phase. These features imply that the fate of electroweak SCPV can in principle be determined with present and near-future experimental data, despite the huge parameter space. The purpose of this Letter is to begin this endeavor, using the most recent results from precision observables and collider searches.

We have found restricted ranges for Higgs masses and the vacuum angles, cf. Eqs. (5)–(7), and have derived a lower bound on charged-Higgs couplings to bottom quarks, cf. Eq. (19). While the remaining parameter space is still huge, it is compact and will be probed from below by precision experiments like EDM measurements and from above by neutral and charged Higgs searches at colliders.

In particular the interplay of lower limits on charged Higgs couplings and upper limits on the charged Higgs mass leads to large production cross sections and branching ratios in channels that have not been explored yet. Our framework thus provides a strong motivation for nonstandard searches at hadron colliders that feature cb or ub associated charged Higgs production and/or decay. We have provided several relevant benchmark points, cf. Table I, which hopefully stimulate more detailed collider studies of these interesting signals that might play an important role in casting the final verdict on the origin of CP violation in weak interactions.

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*ulrich.nierste@kit.edu
†mustafa.tabet@kit.edu
‡robert.ziegler@kit.edu

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