Letter

Entanglement suppression, enhanced symmetry, and a standard-model-like Higgs boson

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We study information-theoretic properties of scalar models containing two Higgs doublets Φ_a, where $a = 1, 2$ is the flavor quantum number. Considering the 2-to-2 scattering $\Phi_a\Phi_b \to \Phi_c\Phi_d$ as a two-qubit system in the flavor subspace and the S-matrix as a quantum logic gate, we analyze the entanglement power of the S-matrix at the tree level, in the limit the gauge coupling is turned off. Demanding the suppression of flavor entanglement during the scattering, the perturbative S-matrix in the broken phase can only be in the equivalent class of the Identity gate and the scalar potential exhibits a maximally enhanced $SO(8)$ symmetry acting on the eight real components of the two doublets. The $SO(8)$ symmetry leads to the alignment limit naturally, giving rise to a Standard-Model-like Higgs boson as a consequence of entanglement suppression.

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Introduction. The concept of symmetry is among the most powerful organizing principles in nature. However, very little has been said about its origin and whether symmetry can be derived from more fundamental principles. On the other hand, J. A. Wheeler famously coined the phrase *It-from-bit*, stipulating that all things physical are information theoretic in its origin [1]. Given that entanglement is one of the most prominent features of quantum mechanics, one wonders if symmetry could arise out of a quantum information-theoretic origin.

Indeed, recent studies in low-energy QCD revealed intriguing connections between the presence of emergent global symmetries and the suppression of spin-entanglement in nonrelativistic scattering of spin-1/2 baryons [2–4]. Of particular interest is the interaction of neutron (n) and proton (p) in the low-energy, which exhibits an approximate $SU(4)_{sm}$ spin-flavor symmetry first observed by E. P. Wigner [5] more than half a century ago, and studied in modern perspective in Ref. [6]. Moreover, s-wave scattering lengths of np are unusually large in both the spin-singlet ($^1S_0$) and spin-triplet ($^3S_1$) channels, which are indicative of nonrelativistic conformal invariance, also known as the Schrödinger symmetry [7,8]. These emergent symmetries in the infrared are usually characterized as fine-tuned or accidental.

Reference [2] first made the fascinating observation that, within the pionless effective field theory [9,10], regions of parameter space where $SU(4)_{sm}$ and Schrödinger symmetry emerge coincide with regions where the spin-entanglement is suppressed in np scattering. In addition, entanglement suppression in flavor-diagonal scattering of octet baryons leads to an even larger $SU(16)$ spin-flavor symmetry [2]. Reference [3] studied these findings in an information-theoretic context and identified the association of the Identity gate with $SU(4)_{sm}$ and $SU(16)$, as well as the SWAP gate with the Schrödinger symmetry. These turned out to be the only two minimal entanglers for two qubit-systems [3]. Subsequently, Ref. [4] extended the analysis to flavor-changing scattering of octet baryons and identified scattering channels whose entanglement suppression are indicative of emergent $SU(6)$, $SO(8)$, $SU(8)$ and $SU(16)$ symmetries.

Given the nascent nature of this subject, it is important to proceed in an exploratory spirit and search for more examples of physical systems exhibiting a correlation between emergent symmetry and entanglement suppression. In this work we study a system of very different nature from the nonrelativistic np interaction; a model of

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electroweak symmetry breaking containing two Higgs doublets $\Phi_a = (\Phi_a^+, \Phi_a^0)^T$, $a = 1, 2$, commonly referred to as two-Higgs-doublet models (2HDMs), which are the prototypical example for physics-beyond-the-Standard Model. We will analyze 2HDMs from a new perspective, focusing on entanglement property of the S-matrix for the scattering $\Phi_a\Phi_b \to \Phi_c\Phi_d$, in the limit the gauge coupling is turned off. The Yukawa coupling, on the other hand, allows us to define flavor quantum number of the Higgs doublets and does not contribute to tree amplitudes. We find that, in the broken phase, requiring the *perturbative* S-matrix to be a minimal entangler in the flavor space leads to a maximal $SO(8)$ symmetry, acting on the eight real components of the two doublets.

The $SO(8)$ symmetry has an important consequence phenomenologically. Since measurements at the Large Hadron Collider indicate properties of the 125 GeV Higgs is standard-model (SM) like [11,12], any viable 2HDMs must be in the "alignment limit" [13–17], where one of the CP-even mass eigenstates is SM-like. It turns out that imposing the $SO(8)$ symmetry, which is broken down to $SO(5)$ upon switching on the $SU(2)_L$ gauge interactions, leads to "natural alignment" [18,19]. Therefore, entanglement suppression in 2HDMs gives rise to a SM-like Higgs boson.

Two-Qubit System Essentials. Here we briefly summarize the key concepts of quantum information needed for the present work. More comprehensive details can be found in Refs. [3,4].

We start with two distinguishable qubits, Alice (A) and Bob (B), each with its own basis vectors $\{|1\rangle_I, |2\rangle_I\}$, $I = A, B$. It is conventional to define the computation basis $\{|11\rangle, |12\rangle, |21\rangle, |22\rangle\}$, where $|ij\rangle = |i\rangle_A \otimes |j\rangle_B$. There are several quantitative measures of entanglement in two-qubit systems [20], although Ref. [3] showed that all of them are related to the concurrence Δ [21,22], which for a normalized state $|\Psi\rangle = c_{11}|11\rangle + c_{21}|21\rangle + c_{12}|12\rangle + c_{22}|22\rangle$ is defined as

$$\Delta(\Psi) = 2|c_{11}c_{22} - c_{12}c_{21}|. \tag{1}$$

The concurrence has a minimum at 0, if $|\Psi\rangle$ is not entangled, and a maximum at 1 if it is maximally entangled. Other commonly employed entanglement measures include the von Neumann entropy $E_{vN}(\rho) = -\mathrm{Tr}(\rho_A \ln \rho_A)$ and the linear entropy $E_L(\rho) = -\mathrm{Tr}[\rho_A(\rho_A - 1)]$, where $\rho = |\Psi\rangle\langle\Psi|$ is the density matrix and $\rho_{A/B} = \mathrm{Tr}_{B/A}(\rho)$ is the reduced density matrix for Alice/Bob.

Entanglement is a property of quantum states. Nevertheless we are more interested in the ability of a quantum operator U to generate entanglement. In this regard, the entanglement power of a unitary operator is defined by averaging over all direct product states that U acts upon [23,24]:

$$\Delta(U) = \overline{\Delta(U|\psi_A > \otimes |\psi_B >)}, \tag{2}$$

where the average is over each Bloch sphere. Importantly, local operators which can be written as the product of singlet-qubit quantum gates, $V = U_A \otimes U_B$ do not generate entanglement. This defines an equivalent class among the two-qubit gates,

$$U \sim U', \quad \text{if } U = V_1 U' V_2. \tag{3}$$

Operators in the same equivalent class have the same entanglement power. Classification of all nonlocal, and hence entanglement generating operators in a two-qubit system has been achieved long ago [20,25,26]. However, for our purpose we focus on entanglement suppressing operators characterized by $\Delta = 0$, which consist of only the Identity gate and the SWAP gate [3], as well as their equivalent classes. In the computational basis they are defined by $\mathbf{1}|ij\rangle = |ij\rangle$ and $\mathrm{SWAP}|ij\rangle = |ji\rangle$. We will represent equivalent classes of 1 and SWAP by $[\mathbf{1}]$ and [SWAP], respectively.

In low-energy QCD, nonrelativistic np scattering is dominated by the s-wave and the S-matrix can be written as [3]

$$S = \frac{1}{2}(e^{2i\delta_1} + e^{2i\delta_0})\mathbf{1} + \frac{1}{2}(e^{2i\delta_1} - e^{2i\delta_0})\ \mathrm{SWAP}, \tag{4}$$

where δ_0 and δ_1 are the scattering phases in the ${}^1S_0$ and ${}^3S_1$ channels, respectively. Then, Alice and Bob are the spin-1/2 proton and neutron, respectively. One can see from Eq. (4) that $S \propto \mathbf{1}$ if $\delta_0 = \delta_1$ and $S \propto \mathrm{SWAP}$ if $|\delta_0 - \delta_1| = \pi/2$. The observation in Ref. [2] is that $\delta_0 = \delta_1$ corresponds to Wigner's $SU(4)_{sm}$ spin-flavor symmetry [5,6] and $|\delta_0 - \delta_1| = \pi/2$ gives rise to the Schrödinger symmetry [7,8]. Both are emergent symmetries not present in the fundamental QCD Lagrangian.

2HDM Essentials. In 2HDM there are two hypercharge-one, $SU(2)$ doublet fields $\Phi_a = (\Phi_a^+, \Phi_a^0)^T$, $a = 1, 2$, and the most general potential is given by, following the notation of Ref. [27],

$$\begin{aligned}\mathcal{V} = {}& m_1^2\Phi_1^\dagger\Phi_1 + m_2^2\Phi_2^\dagger\Phi_2 - [m_{12}^2\Phi_1^\dagger\Phi_2 + \mathrm{H.c.}] \\ &+ \frac{\lambda_1}{2}(\Phi_1^\dagger\Phi_1)^2 + \frac{\lambda_2}{2}(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ &+ \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) \\ &+ \left[\frac{\lambda_5}{2}(\Phi_1^\dagger\Phi_2)^2 + \lambda_6(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2)\right. \\ &\left. + \lambda_7(\Phi_2^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) + \mathrm{H.c.}\right].\end{aligned} \tag{5}$$

For simplicity we assume CP conservation and λ_i are real parameters, although our results can be easily generalized to the CP-violating case. We also assume a $U(1)_{em}$-preserving vacuum, leading to two scalar vacuum expectation values (VEVs), v_1 and v_2, that are real and nonnegative, with $(v_1^2+v_2^2)^{1/2}\equiv v=246$ GeV. We define $t_\beta=v_2/v_1\geq 0$, $0\leq\beta\leq\pi/2$, such that $c_\beta\equiv\cos\beta=v_1/v$ and $s_\beta\equiv\sin\beta=v_2/v$.

Before considering the couplings of the two Higgs doublets to fermions, the flavor quantum number of Higgs doublets is not well-defined. This is because Φ_1 and Φ_2 have identical SM quantum numbers and one is free to redefine the scalar fields by a global $U(2)$ rotation of $\vec{\Phi}=(\Phi_1,\Phi_2)^T$, which leaves the scalar kinetic term invariant, $\vec{\Phi}\to\vec{\Phi}'=\mathcal{U}\vec{\Phi}$, $\mathcal{U}^\dagger\mathcal{U}=\mathbb{I}$. Parameters appearing in Eq. (5) are not invariant under $U(2)$ rotations, whereas the potentials related by $U(2)$ rotations are physically equivalent. One can remove the $U(2)$ redundancy by introducing couplings to fermions. That is, once Yukawa couplings are introduced, flavor can be defined. For example, in type II 2HDMs [28,29], one doublet couples to the up-type fermions while the other couples to the down-type fermions, thereby allowing us to distinguish the two doublets. Another choice of basis, which is convenient for studying the phenomenological property of 2HDMs, [14–17], is the Higgs basis [30], defined by (H_1,H_2) with the property; $\langle H_1^0\rangle=v/\sqrt{2}$ and $\langle H_2^0\rangle=0$. In the Higgs basis the scalar potential is the same as in Eq. (5) but with the coefficients $\{m_1^2,m_2^2,m_{12}^2\}\to\{Y_1,Y_2,Y_3\}$ and $\lambda_i\to Z_i$. The minimization of scalar potential gives $Y_1=-Z_1v^2/2$ and $Y_3=-Z_6v^2/2$. The entanglement power of the S-matrix does not depend on the operator basis since making a $U(2)$ rotation corresponds to single-qubit operation and preserves the entanglement power.

The alignment limit is defined as the case when the scalar $h\equiv\mathrm{Re}(H_1^0)$ coincides with the 125 GeV mass eigenstate. In this case h, which carries the full VEV, couples to the massive gauge bosons with the SM strength when the gauge coupling is turned on. It is shown in Refs. [14–17] that the alignment is achieved by the condition,

$$Z_6=0. \tag{6}$$

In the case of CP-violation, Z_6 is complex and the alignment condition is really two equations; $\mathrm{Re}(Z_6)=0$ and $\mathrm{Im}(Z_6)=0$, which eliminate mass mixings of h with the other two neutral scalars. In any case, when $|Z_6|\ll 1$, h is approximately aligned with the 125 GeV mass eigenstate, which then becomes SM-like. Moreover, in this limit Z_1v^2 is the dominant contribution to the mass of h; $M_h^2\approx Z_1v^2=-2Y_1$. To summarize, the mass of the SM-like Higgs boson is controlled by Z_1 while the departure from Higgs alignment is given by Z_6.

S-matrix as an Identity gate. We now investigate the information-theoretic properties of 2HDMs, focusing on the S-matrix as an entanglement operator in the flavor-space in the scattering $\Phi_a\Phi_b\to\Phi_c\Phi_d$. In terms of Alice and Bob qubits, we identify $|i\rangle_A=\Phi_i^+$ and $|i\rangle_B=\Phi_i^0$, $i=1,2$, respectively. The reason for choosing different electroweak quantum numbers is that Alice and Bob are then associated with distinguishable qubits. The S-matrix, being a unitary operator, then can be thought of as a two-qubit quantum logic gate. Recall that the S-matrix is related to the transition matrix T:

$$S=1+iT, \tag{7}$$

where the matrix elements of the T-matrix are given by

$$\langle\Phi_c\Phi_d|iT|\Phi_a\Phi_b\rangle=i(2\pi)^4\delta^{(4)}(p_a+p_b-p_c-p_c)M_{ab,cd}. \tag{8}$$

$M_{ab,cd}$ are the scattering amplitudes one typically computes in perturbation theory. Notice that the T-matrix, and therefore the amplitude itself, is not a unitary operator and does not admit an interpretation as a quantum gate. In fact, unitarity of the S-matrix requires

$$i(T^\dagger-T)=TT^\dagger, \tag{9}$$

which is nothing but the optical theorem. At the tree-level, the amplitude does not have an imaginary part and the T-matrix is Hermitian. This can be seen from the fact that, if $T\sim\mathcal{O}(\lambda)$ in perturbation, $T^\dagger T\sim\mathcal{O}(\lambda^2)$ is higher order in the coupling constants and the right-hand side of Eq. (9) can be ignored; perturbative unitarity of the S-matrix is fulfilled at $\mathcal{O}(\lambda)$. It is worth pointing out that our approach is different from some in the literature which looked at the entanglement property of the amplitude, instead of the S-matrix [31–33].

We are interested in an S-matrix which suppresses flavor entanglement in 2-to-2 scattering, when turning off the gauge fields. *A priori* we need to consider the two equivalent classes associated with the Identity and the SWAP gates, [$\mathbf{1}$] and [SWAP], respectively [3]. However, we argue that perturbatively the S-matrix could only be in [$\mathbf{1}$] and not [SWAP]. This is most clear by looking at Eq. (7), which implies

$$S\sim[\mathbf{1}]\Leftrightarrow T\sim[\mathbf{1}], \tag{10}$$

$$S\sim[\mathrm{SWAP}]\Leftrightarrow T\sim i([\mathbf{1}]+[\mathrm{SWAP}]). \tag{11}$$

In other words, the S-matrix being in [SWAP] requires a tree-level cancellation between the T-matrix, which we compute in perturbation, against the noninteracting part of the S-matrix. This can only be achieved in a strongly-coupled theory. Indeed, Refs. [7,8] found the SWAP gate

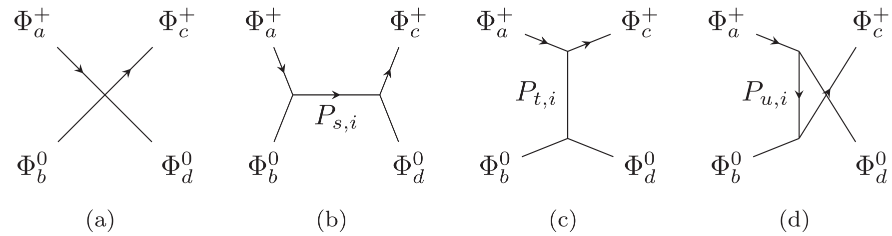

FIG. 1. Feynman diagrams of $\Phi_a^+\Phi_b^0 \to \Phi_c^+\Phi_d^0$ scattering in the symmetry broken phase.

is associated with fermionic systems interacting with the largest strength allowed by unitarity—fermions at unitarity—and Schrödinger symmetry emerges from it. For weakly coupled theories, an entanglement-suppressing S-matrix can only be in [**1**] at finite orders in perturbation theory, except when a certain class of diagrams is resummed to all orders [34].

In what follows we will focus on the flavor subspace of the amplitude $M_{ab,cd}$ defined in Eq. (8), which is Hermitian at the tree level, and work out the conditions on the amplitude in order for the S-matrix to be in [**1**]. Starting from an initial product state in the flavor space, $|\Phi_a\Phi_b\rangle = (\kappa|1\rangle + \epsilon|2\rangle) \otimes (\gamma|1\rangle + \delta|2\rangle)$, where $|\kappa|^2 + |\epsilon|^2 = |\gamma|^2 + |\delta|^2 = 1$. The outgoing state then has the flavor structure,

$$|\Phi_c\Phi_d\rangle = (\delta_{ac}\delta_{bd} + iM_{ab,cd})|\Phi_a\rangle \otimes |\Phi_b\rangle. \tag{12}$$

Demanding that the concurrence in Eq. (1) vanishes, $\Delta(|\Phi_c\Phi_d\rangle) = 0$, and working to the first order in perturbation by keeping only terms linear in $M_{ab,cd}$, we obtain

$$M_{11,11} + M_{22,22} = M_{12,12} + M_{21,21}, \tag{13}$$

$$M_{11,22} = M_{12,21} = M_{21,12} = M_{22,11} = 0, \tag{14}$$

$$\begin{aligned} M_{11,12} &= M_{21,22}, & M_{11,21} &= M_{12,22}, \\ M_{22,21} &= M_{12,11}, & M_{22,12} &= M_{21,11}. \end{aligned} \tag{15}$$

More details can be found in the Supplemental Materials. These are the conditions the tree-level amplitude must satisfy in order for the S-matrix to be in the equivalent class of the Identity gate, $S = [\mathbf{1}]$, which are more general than simply requiring $M_{ab,cd} = \mathbf{1}$. This situation is markedly different from that in the np scattering, where rotational invariance constrains the s-wave S-matrix to be exactly $\mathbf{1}$ perturbatively. If we had imposed the $SU(2)$ flavor symmetry in our 2HDMs, we would have arrived at the same situation.

$SO(8)$ Symmetry. In this section we compute the tree-level scattering amplitude for $\Phi_a^+\Phi_b^0 \to \Phi_c^+\Phi_d^0$ in the broken phase. The goal is to demonstrate that, when the 2-to-2 amplitude minimizes entanglement and satisfies Eqs. (13)–(15), a maximal $SO(8)$ symmetry arises. The 2-to-2 amplitude includes four Feynman diagrams shown in Fig. 1; the 4-point contact interaction and the $s/t/u$-channels mediated by cubic vertices in the broken phase. The internal propagators in Fig. 1 necessitates a rotation into the mass eigenstates, which in general is different between the charged sector and the neutral sector. However, an advantage of the Higgs basis is that the charged sector is already diagonal since $t_\beta = 0$. So we will perform that calculation in the Higgs basis, $H_1 = (G^+, v/\sqrt{2} + H_1^0)^T$ and $H_2 = (H^+, H_2^0)^T$, where G^+ is the charged Goldstone boson and H^+ is the massive charged scalar. In the neutral sector there are four mass eigenstates which we denote by (h, H, G^0, A); h is the lightest CP-even scalar, which we assume to be the 125 GeV Higgs boson, H and $A = \mathrm{Im}[H_2^0]$ are the CP-even and CP-odd heavy scalars, respectively, and $G^0 = \mathrm{Im}[H_1^0]$ is a Goldstone boson. The rotation matrix $\mathcal{R}$ in the neutral sector is given by

$$\begin{pmatrix} h \\ H \\ G^0 \\ A \end{pmatrix} = \mathcal{R}\begin{pmatrix} H_1^0 \\ H_1^{0*} \\ H_2^0 \\ H_2^{0*} \end{pmatrix}, \quad \mathcal{R} = \frac{1}{2}\begin{pmatrix} -s_{\tilde\alpha} & -s_{\tilde\alpha} & c_{\tilde\alpha} & c_{\tilde\alpha} \\ c_{\tilde\alpha} & c_{\tilde\alpha} & s_{\tilde\alpha} & s_{\tilde\alpha} \\ -i & i & 0 & 0 \\ 0 & 0 & -i & i \end{pmatrix}, \tag{16}$$

where $\tilde\alpha$ is the mixing angle in the neutral CP-even sector in the Higgs basis. It is related to the corresponding mixing angle α in the general basis by $\tilde\alpha = \alpha - \beta$. Observe that the alignment condition corresponds to $c_{\tilde\alpha} = 0$.

The full amplitude is

$$iM_{ab,cd} = iM^0_{ab,cd} - \frac{v^2}{2}\sum_i \sum_{r=s,t,u} M^r_{i\ ab,cd} P_{r,i}, \tag{17}$$

$$M^0_{ab,cd} = \begin{pmatrix} Z_1 & Z_6 & Z_6 & Z_5 \\ Z_6 & Z_3 & Z_4 & Z_7 \\ Z_6 & Z_4 & Z_3 & Z_7 \\ Z_5 & Z_7 & Z_7 & Z_2 \end{pmatrix}, \tag{18}$$

$$M^s_{i\ ab,cd} = M_{abi}M^*_{cdi}, \qquad M^u_{iab,cd} = M_{adi}M^*_{cbi}, \tag{19}$$

$$M^t_{i\ ab,cd} = \sum_{j,k} \mathcal{R}_{ij} M_{ajc}(\mathcal{R}_{ik} M_{dkb,0})^* + \text{H.c.}, \tag{20}$$

where the propagators entering the $s/t/u$-channel diagrams are $P_{t,i} = i/(t - m_{0,i}^2)$ and $P_{r,i} = i/(r - m_{+,i}^2)$, for $r = s$, u. Masses in the propagator run through $m_{0,i} = \{m_h, m_H, 0, m_A\}$ and $m_{+,i} = \{m_{H^\pm}, 0\}$. Moreover, the cubic vertices M_{dkb} and $M_{dkb,0}$ are

$$\left.\frac{\partial \mathcal{V}}{\partial v}\right|_{v=0} = \frac{1}{\sqrt{2}} \sum_{a,b,c} \left[M_{abc} H_a^+ H_b^0 H_c^- + \frac{1}{2} M_{abc,0} H_a^0 H_b^0 H_c^{0\,*} + \text{H.c.} \right]. \quad (21)$$

In order for the S-matrix to minimize entanglement and be in [**1**] for arbitrary kinematics, we will demand that every term in Eq. (17) satisfies the conditions in Eqs. (13)–(15). For $M^0_{ab,cd}$ in Eq. (18), they lead to $Z_1 + Z_2 = 2Z_3$, $Z_4 = Z_5 = 0$, and $Z_6 = Z_7$. These relations greatly simplify expressions in $M^r_{i\ ab,cd}$. Solving for entanglement suppressing amplitudes in the $s/t/u$-channel then requires [35]:

$$Z_1 = Z_2 = Z_3 \equiv Z, \qquad Z_i = 0, \quad i \neq 1,2,3, \quad (22)$$

$$Y_1 = Y_2 \equiv Y = -Zv^2/2, \qquad Y_3 = 0, \quad (23)$$

which lead to the scalar potential,

$$\mathcal{V} = Y(H_1^\dagger H_1 + H_2^\dagger H_2) + \frac{Z}{2}(H_1^\dagger H_1 + H_2^\dagger H_2)^2 = \frac{Z}{2}\left(\sum\nolimits_{i=1,2} |H_i^0|^2 + G^+G^- + H^+H^- - \frac{v^2}{2}\right)^2. \quad (24)$$

(Complete Feynman rules and full expressions for the amplitudes can be found in the Supplementary Material.) The above potential exhibits a maximal $SO(8)$ symmetry acting on the eight real components of the two doublets and is spontaneously broken to $SO(7)$. The spectrum contains a massive scalar h with $m_h^2 = -2Y = Zv^2$, while all other scalars are exact Goldstone bosons and massless. However, recall that the $SO(8)$ symmetry is explicitly broken by Yukawa and gauge couplings (when turned on) and the Goldstone bosons will either become massive or be "eaten" by the W and Z bosons. Furthermore, to achieve a realistic mass spectrum consistent with null searches at the LHC, $SO(8)$ needs to be broken softly by the mass terms [19]. Since one of the minimization conditions relates Y_3 to Z_6, which controls the alignment condition, one could leave $Y_3 = 0$ and introduce an additional Y_2 contribution, which fixes the nonstandard Higgs spectrum $m_H^2 = m_A^2 = m_{H^\pm}^2 = Y_2 + Zv^2/2$ (see, for example, Ref. [36]). The latter clearly shows that, in the $SO(8)$ symmetric limit, Eq. (23) leads to massless nonstandard Higgs bosons.

Conclusions. In this work we analyzed information-theoretic properties of general 2HDMs, a prototypical example for beyond-the-SM theories. When the gauge and Yukawa couplings are turned off, demanding that the perturbative S-matrix suppresses flavor entanglement in $\Phi_a\Phi_b \to \Phi_c\Phi_d$, and is in the equivalent class of the Identity gate, singles out regions of parameter space where the $SU(2) \times U(1)$ symmetry is enhanced to $SO(8)$, which in turn is broken spontaneously to $SO(7)$ and gives rise to a SM-like Higgs boson as a consequence of entanglement suppression. Turning on the Yukawa and gauge couplings results in explicit $SO(8)$ breaking and lifts the otherwise massless non-SM Higgs bosons. However, at the tree-level Yukawas do not enter the 2-to-2 scattering of scalars and therefore will not affect the conditions for entanglement suppression. As for the impact of gauge couplings on entanglement suppression, we plan to investigate in the near future. In any case, a realistic spectrum compatible with current LHC bounds requires the $SO(8)$ to be further broken softly by mass terms and the entanglement suppression is approximate. We leave for future work a detailed analysis of the degree of entanglement suppression and its phenomenological implications for LHC data. In particular, there are several symmetry groups leading to the alignment conditions which are discovered in Refs. [18,19], but only a specific symmetry group [for instance the $SO(8)$ in the absence of gauge and Yukawa couplings] is predicted by the entanglement suppression. Therefore, an experimental test of the predicted symmetry would be an experimental test of entanglement suppression, but a test of the other symmetries will not be related to entanglement properties.

To summarize, the unexpected connections between quantum entanglement, emerging symmetries and a SM-like Higgs boson provide potentially rich and fruitful insights for the exploration of physics beyond-the-SM using tools in quantum information science.

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