

## Does the HyperCP Evidence for the Decay $\Sigma^+ \rightarrow p\mu^+\mu^-$ Indicate a Light Pseudoscalar Higgs Boson?

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The HyperCP Collaboration has observed three events for the decay  $\Sigma^+ \rightarrow p\mu^+\mu^-$  which may be interpreted as a new particle of mass 214.3 MeV. However, existing data from kaon and  $B$ -meson decays provide stringent constraints on the construction of models that support this interpretation. In this Letter we show that the “HyperCP particle” can be identified with the light pseudoscalar Higgs boson in the next-to-minimal supersymmetric standard model, the  $A_1^0$ . In this model there are regions of parameter space where the  $A_1^0$  can satisfy all the existing constraints from kaon and  $B$ -meson decays and mediate  $\Sigma^+ \rightarrow p\mu^+\mu^-$  at a level consistent with the HyperCP observation.

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Three events for the decay mode  $\Sigma^+ \rightarrow p\mu^+\mu^-$  with a dimuon invariant mass of 214.3 MeV have been recently observed by the HyperCP Collaboration [1]. It is possible to account for these events within the standard model (SM) [2], but the probability of having all three events at the same dimuon mass, given the SM predictions, is less than 1%. This suggests a new-particle interpretation for these events, for which the branching ratio is  $(3.1_{-1.9}^{+2.4} \pm 1.5) \times 10^{-8}$  [1].

The existence of a new particle with such a low mass would be remarkable as it would signal the existence of physics beyond the SM unambiguously. It would also be very surprising because this low-energy region has been thoroughly explored by earlier experiments studying kaon and  $B$ -meson decays. The challenge posed by a new-particle interpretation of the HyperCP events is therefore manifold. It requires a new-physics model containing a suitable candidate for the new particle,  $X$ , which explains why it is light. It also requires an explanation of why  $X$  has not been observed by other experiments that covered the same kinematic range. Finally, it requires that the interactions of  $X$  produce the rate implied by the HyperCP observation.

In this Letter we show that there is a model, the next-to-minimal supersymmetric standard model (NMSSM) [3], containing a light pseudoscalar Higgs particle that can satisfy all existing constraints and is therefore a candidate explanation for the HyperCP events. The model contains more than one Higgs particle, and it is the lightest one, the  $A_1^0$ , that can be identified with  $X$ .

The possibility that  $X$  mediated the HyperCP events has been explored to some extent in the literature [4–6], where it has been shown that kaon decays place severe constraints on the flavor-changing two-quark couplings of  $X$ . It has

also been shown [7] that a light sgoldstino is a viable candidate for  $X$ . It is well known in the case of light Higgs boson production in kaon decay that, in addition to the two-quark flavor-changing couplings, there are comparable four-quark contributions [8]. They arise from the combined effects of the usual SM four-quark  $[\Delta S] = 1$  operators and the flavor-conserving couplings of  $X$ . We have recently computed the analogous four-quark contributions to light Higgs production in hyperon decay [9] and found that they can also be comparable to the two-quark contributions previously discussed in the literature.

The interplay between the two- and four-quark contributions makes it possible to find models with a light Higgs boson responsible for the HyperCP events that has not been observed in kaon or  $B$ -meson decay. However, it is not easy to devise such models respecting all the experimental constraints. In most models that can generate  $\bar{d}sX$  couplings, the two-quark operators have the structure  $\bar{d}(1 \pm \gamma_5)sX$ . Since the part without  $\gamma_5$  contributes significantly to  $K \rightarrow \pi\mu^+\mu^-$ , their data imply that these couplings are too small to account for the HyperCP events [4–6]. In some models, there may be parameter space where the four-quark contributions mentioned above and the two-quark ones are comparable and cancel sufficiently to lead to suppressed  $K \rightarrow \pi\mu^+\mu^-$  rates while yielding  $\Sigma^+ \rightarrow p\mu^+\mu^-$  rates within the required bounds. However, since in many models the flavor-changing two-quark couplings  $\bar{q}q'X$  are related for different  $(q, q')$  sets, experimental data on  $B$ -meson decays, in particular,  $B \rightarrow X_s\mu^+\mu^-$ , also provide stringent constraints. For these reasons, the light (pseudo)scalars in many well-known models, such as the SM and the two-Higgs-doublet model, are ruled out as candidates to explain the HyperCP events [9].

In this Letter we show that the  $A_1^0$  in the NMSSM can satisfy the constraints required of  $X$ , namely,

$$\begin{aligned}\mathcal{B}(K^\pm \rightarrow \pi^\pm A_1^0) &\lesssim 8.7 \times 10^{-9}, \\ \mathcal{B}(K_S \rightarrow \pi^0 A_1^0) &\lesssim 1.8 \times 10^{-9}, \\ \mathcal{B}(B \rightarrow X_s A_1^0) &\lesssim 8.0 \times 10^{-7},\end{aligned}\quad (1)$$

obtained in Ref. [9] from the measurements of  $K^\pm \rightarrow \pi^\pm \mu^+ \mu^-$  [10,11],  $K_S \rightarrow \pi^0 \mu^+ \mu^-$  [12], and  $B \rightarrow X_s \mu^+ \mu^-$  [13,14]. We include both two- and four-quark contributions to the  $K$  and  $\Sigma^+$  decays, neglecting the latter contributions to the  $B$  decay.

In the NMSSM there is a gauge-singlet Higgs field  $N$  in addition to the two Higgs fields  $H_u$  and  $H_d$  responsible for the up- and down-type quark masses in the minimal supersymmetric standard model. We follow the specific model described in Ref. [15], with suitable modifications. The superpotential of the model is given by

$$W = QY_u H_u U + QY_d H_d D + LY_e H_d E + \lambda H_d H_u N - \frac{1}{3}kN^3, \quad (2)$$

where  $Q$ ,  $U$ ,  $D$ ,  $L$ , and  $E$  represent the usual quark and lepton fields,  $Y_{u,d,e}$  are the Yukawa couplings, and  $\lambda$  and  $k$  are dimensionless parameters. The soft-supersymmetry-breaking term in the Higgs potential is

$$\begin{aligned}V_{\text{soft}} &= m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_N^2 |N|^2 \\ &\quad - (\lambda A_\lambda H_d H_u N + \frac{1}{3}k A_k N^3 + \text{H.c.}),\end{aligned}\quad (3)$$

and the resulting Higgs potential has a global  $U(1)_R$  symmetry in the limit that the parameters  $A_\lambda, A_k \rightarrow 0$  [16].

There are two physical pseudoscalar bosons in the model which are linear combinations of the pseudoscalar components in  $H_{u,d}$  and  $N$ . The lighter one,  $A_1^0$ , has a mass given by

$$m_{\mathcal{A}}^2 = 3kx A_k + \mathcal{O}(1/\tan\beta) \quad (4)$$

in the large- $\tan\beta$  limit, where  $x = \langle N \rangle$  is the vacuum expectation value of  $N$  and  $\tan\beta$  is the ratio of vacuum expectation values of the two Higgs doublets. If the  $U(1)_R$  symmetry is broken slightly, the mass of  $A_1^0$  becomes naturally small, with values as low as  $\sim 100$  MeV phenomenologically allowed [15,16]. We will now show that this  $CP$ -odd scalar can play the role of the  $X$  particle.

At tree level, the  $A_1^0$  couplings to up- and down-type quarks and to leptons in this model are diagonal and described by

$$\begin{aligned}\mathcal{L}_{\mathcal{A}qq} &= -(l_u m_u \bar{u} \gamma_5 u + l_d m_d \bar{d} \gamma_5 d) \frac{iA_1^0}{v}, \\ \mathcal{L}_{\mathcal{A}\ell} &= \frac{ig_\ell m_\ell}{v} \bar{\ell} \gamma_5 \ell A_1^0,\end{aligned}\quad (5)$$

in the general notation of Ref. [9] used to compute the effect of four-quark operators, where

$$l_d = -g_\ell = \frac{v\delta_-}{\sqrt{2}x}, \quad l_u = \frac{l_d}{\tan^2\beta}, \quad (6)$$

with  $v = 246$  GeV being the electroweak scale and  $\delta_- = (A_\lambda - 2kx)/(A_\lambda + kx)$ . Thus,  $l_u$  can be neglected in the large- $\tan\beta$  limit. The interactions in  $\mathcal{L}_{\mathcal{A}qq}$  combined with the operators due to  $W$  exchange between quarks induce the four-quark contributions discussed earlier.

For an  $A_1^0$  of mass 214.3 MeV as we propose, the decay into a muon-antimuon pair completely dominates over the other kinematically allowed modes:  $A_1^0 \rightarrow e^+ e^-$ ,  $\nu \bar{\nu}$ ,  $\gamma\gamma$ . Accordingly, we assume  $\mathcal{B}(X \rightarrow \mu^+ \mu^-) \sim 1$ . In addition, the muon anomalous magnetic moment imposes the constraint [4]

$$|g_\ell| \lesssim 1.2. \quad (7)$$

This bound results in an  $A_1^0$  width  $\Gamma_{\mathcal{A}} \lesssim 3.7 \times 10^{-7}$  MeV, which is consistent with the estimate of  $10^{-7}$  MeV for the width of the HyperCP particle [6] if  $|g_\ell| = |l_d| = \mathcal{O}(1)$ . An  $l_d$  of order one is also consistent with the general estimates [15] for the size of  $v\delta_-/(\sqrt{2}x)$ .

At one-loop level, the two-quark flavor-changing Lagrangian for  $s \rightarrow dA_1^0$  can be written as

$$\mathcal{L}_{\mathcal{A}sd} = \frac{ig_{\mathcal{A}}}{v} [m_s \bar{d}(1 + \gamma_5)s - m_d \bar{d}(1 - \gamma_5)s] A_1^0 + \text{H.c.}, \quad (8)$$

where

$$g_{\mathcal{A}} = \frac{G_F}{8\sqrt{2}\pi^2} \sum_{q=u,c,t} V_{qd}^* V_{qs} C_{\mathcal{A}q}, \quad (9)$$

with  $V_{qr}$  being the elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix and  $C_{\mathcal{A}q}$  depending on the particles in the loops [15]. The corresponding coupling constant  $g'_{\mathcal{A}}$  in  $b \rightarrow sA_1^0$  has a similar expression in terms of  $V_{qs}^* V_{qb} C_{\mathcal{A}q}$ .

In the large- $\tan\beta$  limit, the pseudoscalar  $A_1^0$  in this model does not couple to up-type quarks and Eq. (9) is completely determined by parameters in the supersymmetry (SUSY) sector. In particular,  $C_{\mathcal{A}q}$  can be zero if a super-Glashow-Iliopoulos-Maiani (super-GIM) mechanism is operative [15]. In this limit, the decays  $\Sigma^+ \rightarrow pA_1^0$  and  $K \rightarrow \pi A_1^0$  are dominated by four-quark contributions [9], and the mode  $b \rightarrow sA_1^0$  is negligible. However, the results of Ref. [9] imply that a vanishing  $g_{\mathcal{A}}$  due to  $C_{\mathcal{A}q} \rightarrow 0$  would not satisfy all the existing constraints.

This leads us to some details of the model. Assuming that there is only top squark mixing, and that sup and scharm are degenerate, we have for  $s \rightarrow dA_1^0$  or  $b \rightarrow sA_1^0$  [15],

$$\begin{aligned}C_{\mathcal{A}q} &= \delta_{qt} \tan\beta \left[ -l_d \sum_{i=1,2} X_i \right. \\ &\quad \left. + \lambda m_t v \sin\theta_i \cos\theta_i \sum_{u,w=1,2} Y_{uw} \right],\end{aligned}\quad (10)$$

where  $X_i$  and  $Y_{uw}$  are loop-integral functions of parameters in the chargino and squark mixing matrices given in Ref. [15], and  $\theta_{\tilde{t}}$  is the top squark mixing angle. Consequently, in this scenario the CKM factors involving the top quark are the relevant ones. It then follows from Eq. (9) that  $C_{\mathcal{A}q}$  of the right magnitude for  $\mathcal{B}(\Sigma^+ \rightarrow pA_1^0) \sim 3 \times 10^{-8}$  will yield  $\mathcal{B}(B \rightarrow X_s A_1^0)$  that violates its limit in Eq. (1) by about 3 orders of magnitude.

Now, the penguin amplitude in this model is dominated by squark and chargino intermediate states, and the particular scaling with the top-quark CKM factor in Eq. (10) follows from the assumption of top squark mixing only. If we assume instead that there is only sup mixing,  $C_{\mathcal{A}q}$  will be proportional to  $\delta_{qu}$ . In this case, it is possible for  $A_1^0$  to be the HyperCP particle without violating the bounds in Eq. (1). To be specific, assuming that sup mixing is dominant and the top and charm squarks are degenerate, we obtain

$$C_{\mathcal{A}q} = \delta_{qu} \tan\beta \left[ -l_d \sum_{i=1,2} X_i + \lambda m_u v \sin\theta_{\tilde{u}} \cos\theta_{\tilde{u}} \sum_{u,w=1,2} Y_{uw} \right]. \quad (11)$$

In addition, Eq. (11) has a vanishing  $CP$ -violating phase if all the SUSY parameters are real. This is needed if the contributions from two- and four-quark operators are to partially cancel each other [9], as needed to satisfy the kaon bounds in Eq. (1). To simplify the discussion, we now consider the case where  $\theta_{\tilde{u}} = 0$ . In this case,

$$g_{\mathcal{A}} = \frac{G_F}{8\sqrt{2}\pi^2} V_{ud}^* V_{us} \sqrt{2} m_W \tan\beta \sum_i m_{\chi_i} U_{i2} V_{i1} [D_3(y_{c,i}) - D_3(y_{u,i})] l_d, \quad (12)$$

where  $U_{ij}$  and  $V_{ij}$  are the unitary matrices diagonalizing the chargino mass matrix,  $D_3(y) = y \ln y / (1 - y)$  is a loop-integral function, and  $y_{r,i} = m_{\tilde{r}}^2 / m_{\chi_i}^2$ , with  $m_{\chi_i}$  being a chargino mass [15]. Therefore, depending on whether the sup mass is larger or smaller than the scharm mass,  $g_{\mathcal{A}}$  can take either sign.

The expressions for the four-quark contributions to the kaon and hyperon decays are lengthy and in general depend on the effective coupling of  $A_1^0$  to a gluon pair [9]. For the simple case above, the contributions to  $A_1^0 gg$  are from quarks in the loop in the large- $\tan\beta$  limit, and the squarks do not contribute since there is no left- and right-squark mixing. Since  $l_u \rightarrow 0$ , the four-quark contributions can be written purely in terms of  $l_d$ . Using the results in Ref. [9] (with  $D - F = 0.25$ ), the full amplitudes are

$$i\mathcal{M}(K^\pm \rightarrow \pi^\pm A_1^0) = g_{\mathcal{A}} \frac{m_K^2 - m_\pi^2}{v} - (1.08 \times 10^{-7}) l_d \frac{m_K^2}{v}, \quad (13)$$

$$\mathcal{M}(\Sigma^+ \rightarrow pA_1^0) = \left[ g_{\mathcal{A}} \frac{m_\Sigma - m_p}{v} - (6.96 \times 10^{-7}) l_d \frac{f_\pi}{v} \right] i\bar{p}\Sigma^+ - \left[ (0.25) g_{\mathcal{A}} \frac{m_\Sigma + m_p}{v} \frac{m_K^2}{m_K^2 - m_\pi^2} + (5.72 \times 10^{-6}) l_d \frac{f_\pi}{v} \right] i\bar{p}\gamma_5 \Sigma^+. \quad (14)$$

The corresponding amplitude for  $K_S \rightarrow \pi^0 A_1^0$ , which we have not displayed, can be derived from Ref. [9].

Choosing, for example,  $l_d = 0.35$ , we show the resulting branching ratios in Fig. 1. Even though it is possible for  $A_1^0$  to explain the HyperCP observation for a wide range of parameters, as the figure indicates, we have found that there are only rather narrow ranges of  $g_{\mathcal{A}}$  and  $l_d$  for which the kaon bounds are also satisfied, namely,  $1.0 \times 10^{-7} \leq g_{\mathcal{A}}/l_d \leq 1.3 \times 10^{-7}$ . The  $B$ -meson bound is automatically satisfied with our assumptions above.

We remark that  $g_{\mathcal{A}}$  in Eq. (12) is proportional to  $G_F V_{ud}^* V_{us} \tan\beta m_W m_{\chi_i}$ , which is of  $\mathcal{O}(1)$  for  $\tan\beta = \mathcal{O}(30)$  and  $m_{\chi_i} = \mathcal{O}(100 \text{ GeV})$ . Thus, in order to have  $g_{\mathcal{A}} = \mathcal{O}(10^{-7})$  with these parameter choices, one needs  $U_{i2} V_{i1} [D_3(y_{c,i}) - D_3(y_{u,i})]$  to be of order  $10^{-7}$  as well. Since, as mentioned above, there is a super-GIM mechanism that can be at work, as long as  $m_{\tilde{u}}$  and  $m_{\tilde{c}}$  are degenerate enough, one can get an arbitrarily small  $g_{\mathcal{A}}$  to satisfy the required value. It turns out that there is also SUSY parameter space where the desired value of  $g_{\mathcal{A}}$  can be produced without very fine-tuned degeneracy.

We have carried out a detailed exploration of the parameter space. In general, if the chargino masses are larger than the squark masses, fine-tuned squark masses, with splitting of less than 1 GeV between sup and scharm masses, are required. However, if the chargino masses [ $\mathcal{O}(100 \text{ GeV})$ ] are smaller than the squark masses [ $\mathcal{O}(1 \text{ TeV})$ ], there is some parameter space satisfying the

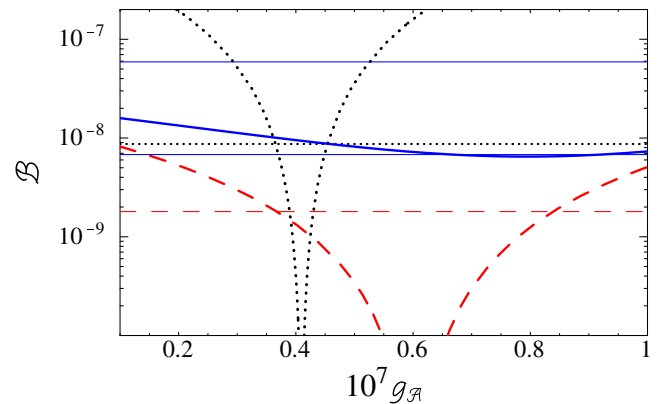


FIG. 1 (color online). Branching ratios of  $\Sigma^+ \rightarrow pA_1^0$  (solid curves),  $K^+ \rightarrow \pi^+ A_1^0$  (dotted curves), and  $K_S \rightarrow \pi^0 A_1^0$  (dashed curves) as functions of  $g_{\mathcal{A}}$  for  $l_d = 0.35$ . The horizontal lines indicate the bounds from the HyperCP result and the kaon bounds from Eq. (1). The allowed range in this case is  $0.38 \leq 10^7 g_{\mathcal{A}} \leq 0.43$ .

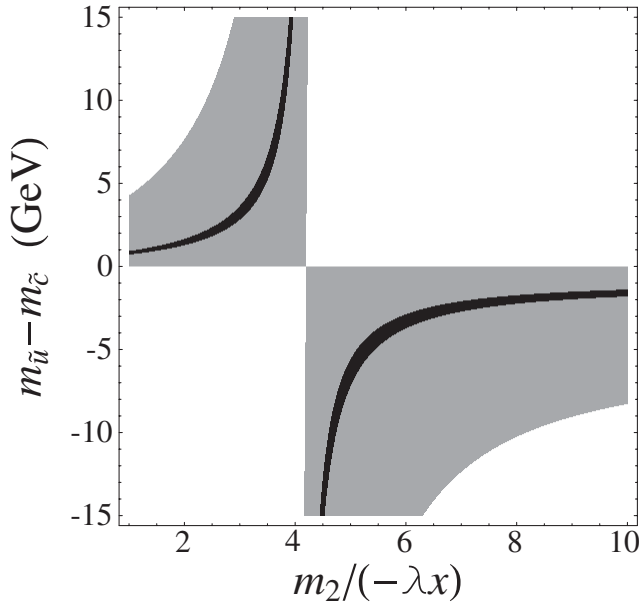


FIG. 2. Parameter space for  $m_{\tilde{u}} - m_{\tilde{c}}$  and  $m_2/(-\lambda x)$  where  $A_1^0$  can explain the HyperCP events (gray regions) and simultaneously satisfy the kaon bounds (black regions), with the input described in the text.

constraints with squark-mass splitting of a few GeV to tens of GeV. Taking  $-\lambda x = 150$  GeV,  $m_{\tilde{c}} = 2.5$  TeV, and  $\tan\beta = 30$ , we display in Fig. 2 the allowed ranges of  $m_{\tilde{u}} - m_{\tilde{c}}$  and  $m_2/(-\lambda x)$  corresponding to the required values  $1.0 \times 10^{-7} \lesssim g_{\mathcal{A}}/l_d \lesssim 1.3 \times 10^{-7}$ , where  $m_2$  is the first element of the chargino undiagonalized mass matrix [15]. In this case, the two chargino eigenmasses are, respectively,  $100 \text{ GeV} \lesssim m_{\chi_1} \lesssim 150 \text{ GeV}$  and  $200 \text{ GeV} \lesssim m_{\chi_2} \lesssim 1500 \text{ GeV}$ . These chargino masses, especially the lighter one, are in the interesting range which may be probed at the LHC and ILC. In this figure, we also display the much larger (gray) regions where  $A_1^0$  can explain the HyperCP observation.

In conclusion, we have shown that a light pseudoscalar Higgs boson in the NMSSM can be identified as the possible new particle responsible for the HyperCP events while satisfying all constraints from kaon and  $B$ -meson decays. We have found that there are large regions in the parameter space where the particle can explain the HyperCP observation. This by itself is not trivial, considering that the SM, for example, yields rates that are too large by more than a factor of 10 [9]. Only after the kaon bounds are imposed does the allowed parameter space

become much more restricted. We believe that this is sufficiently intriguing to warrant a revisiting of these kaon bounds, perhaps by the NA48 experiment.

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