The Pseudo-Nambu Goldstone Boson of Metastable SUSY-Violation

Howard E. Haber

SM and BSM physics at the LHC

CERN Theory Institute

24 August 2009

Based on work in collaboration with Tom Banks (arXiv:0908.2004 [hep-ph])

Outline

- The Pseudo-Nambu-Goldstone boson (PNGB) in dynamical SUSY-breaking models
- C and P quantum numbers of the PNGB
- How does the PNGB couple to Standard Model (SM) gauge bosons?
- Coupling of the PNGB to SM hadronic vector currents
- PNGB Yukawa couplings to SM fermions
- Phenomenological Constraints
- Conclusions

Viable fundamental theories of SUSY-breaking

A viable and compelling fundamental theory of SUSY-breaking (which determines the origin of SUSY-breaking mass parameters that enter as coefficients of relevant operators) is not easily established.

In 2006, Intriligator, Seiberg and Shih (ISS) argued that the space of viable models would be considerably enlarged if one allowed for metastable SUSY-breaking (i.e., local SUSY-breaking minima that are not global). As long as the lifetime of the metastable vacuum is sufficiently long, the corresponding model is a potential candidate for the fundamental SUSY-breaking of our world.

The framework for metastable SUSY-breaking

We shall employ an ISS-type model to provide the fundamental source of SUSY-breaking for the MSSM. This model will consist of an SU(N_c) super Yang Mills theory consisting of N_F flavors of vector-like quarks (denoted henceforth by $\mathcal{Q}$). An SU(3)$\times$SU(2)$\times$U(1) subgroup of the global SU$(N_F)\times$SU(N_F) flavor symmetry is gauged and identified as the Standard Model gauge group. This is a model of *direct* gauge mediation of SUSY-breaking.

In the dual magnetic theory (assuming $N_c < N_F < \frac{3}{2}N_c$), some of the dual quarks acquire non-zero vevs at the SUSY-breaking metastable minimum. The hidden sector quarks carry *meta-baryon number* [“meta” is used to distinguish this from ordinary baryon number carried by the SM quarks], which is spontaneously broken at the ISS scale, denoted by $\Lambda_{\rm ISS}$.*

---

*Since the ISS sector is vector-like with respect to the SU(N_C), the meta-baryon number is a non-anomalous global symmetry.

The PNGB of metastable SUSY-violation

The spontaneous breaking of meta-baryon number yields an exactly massless Goldstone boson, $\mathcal{P}$. But, we do not expect the meta-baryon number global symmetry to be exact to arbitrarily high energy scales.

Let $M_U \gg \Lambda_{\rm ISS}$ be the scale at which the global meta-baryon number is explicitly broken. Taking the (irrelevant) operators generated at this high-scale into account, the Goldstone boson of spontaneous meta-baryon number breaking acquires a mass—it is now a pseudo-Nambu-Goldstone boson (PNGB).

The lowest gauge-invariant operator (involving fields of the electric theory) that violates meta-baryon number is

$$\delta W \sim \frac{1}{\Lambda_U^{N_c-3}} \mathcal{Q}^{N_c} \, ,$$

where we take $N_c > 3$ (otherwise the above operator is no longer irrelevant).

If for some reason, the above operator is disallowed (due, say, to discrete symmetries preserved at the scale M_U) one can introduce an extra singlet field S and choose,

$$\delta W \sim \frac{1}{\Lambda_U^{N_c+p-3}} \mathcal{Q}^{N_c} S^p \, ,$$

for some suitably chosen p. In either case, $\mathcal{P}$ acquires a non-trivial potential due to the explicit breaking:

$$V \sim \Lambda_{\rm ISS}^4 \left(\frac{\Lambda_{\rm ISS}}{M_U} \right)^{N_c+p-3} U(\mathcal{P}/\Lambda_{\rm ISS}) \, ,$$

where $U(x) = U_0 + cx^2 + \dots$. Consequently,

$$\boxed{m_{\mathcal{P}}^2 \sim \Lambda_{\rm ISS}^2 \left(\frac{\Lambda_{\rm ISS}}{M_U} \right)^{N_c+p-3} \, .}$$

A range of possible PNGB masses

The choice of $M_{\rm ISS}$ and M_U is highly model-dependent. In the framework of gauge-mediated SUSY-breaking models, one expects $\Lambda_{\rm ISS}$ to be in the TeV to multi-TeV range. In this work we choose:

$$\Lambda_{\rm ISS} \sim 2 \text{ TeV} ,$$

which is a wildly optimistic choice (most probably this scale is significantly larger).

For the high-energy scale M_U, one can imagine a number of possible choices:

- the reduced Planck scale (2×10^{18} GeV)
- the grand unification scale (2×10^{16} GeV)
- the right-handed neutrino (seesaw) scale (5×10^{14} GeV)

Finally, we choose N_c and p. As we previously indicated that $N_c > 3$, we consider three values $N_c + p = 4, 5, 6$. Taking the extremes yields a range of masses

$$m_{\mathcal{P}} \sim 6 \times 10^{-11}\ \text{eV}\ —\ 4\ \text{MeV} .$$

PNGB quantum numbers

The ISS sector consists of vector-like quarks with SU(N_c) gauge-interactions that conserve C and P separately.

Thus, we can assign definite C and P quantum numbers to $\mathcal{P}$. Since the meta-baryon current is a vector (not axial vector) current, it follows from:

$$\langle 0|J^\mu_{MB}(0)|\mathcal{P}\rangle = f_{\mathcal{P}} q^\mu \,,$$

that $C(\mathcal{P}) = -1$ and $P(\mathcal{P}) = +1$. In contrast, the SM pion, for which $\langle 0|A^\mu(0)|\pi\rangle = f_\pi q^\mu$ (where $A^\mu = \bar{u}\gamma^\mu\gamma_5 d$ is an axial vector current) implies that $C(\pi) = +1$ and $P(\pi) = -1$.

Thus, $\mathcal{P}$ is a CP-odd, C-odd scalar. As such, it cannot couple diagonally to a fermion-antifermion pair or scalar particle antiparticle pair.

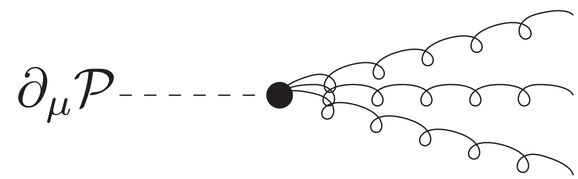
How to couple a CP-odd, C-odd scalar to gauge bosons

By C-invariance, a C-odd scalar can only couple to an odd number of photons. In particular, $\mathcal{P}$ cannot couple to two photons. Likewise, $\mathcal{P}$ cannot couple to two gluons (although it can couple to $2n$ gluons for $n \geq 2$ by making use of the f_{abc} tensor).

The C and P conserving coupling of $\mathcal{P}$ to three photons (or gluons) with the least number of derivatives is unique (Dolgov 1968):

$$\boxed{\mathcal{L}_{\rm eff} = \frac{g_3^3}{\Lambda_{\rm ISS}^6} d_{abc} (D_\rho G_{\alpha\beta})^a (D^\beta G_{\sigma\tau})^b (D^\rho D^\alpha G^{\sigma\tau})^c \mathcal{P}\,,}$$

where $(D_\rho G_{\alpha\beta})^a \equiv D_\rho^{ab} G_{\alpha\beta}^b$, etc. Here, $G_{\mu\nu}^a \equiv \partial^\mu A_\nu^a - \partial^\nu A_\mu^a - g f_{abc} A_\mu^a A_\nu^b$ is the gluon field strength tensor, $D_\mu^{ab} \equiv \delta^{ab}\partial_\mu + g f_{abc} A_\mu^c$ is the covariant derivative acting on an adjoint field, and $d_{abc} \equiv 2{\rm Tr}(\{T_a, T_b\} T_c)$ is the totally symmetric tensor of color SU(3).

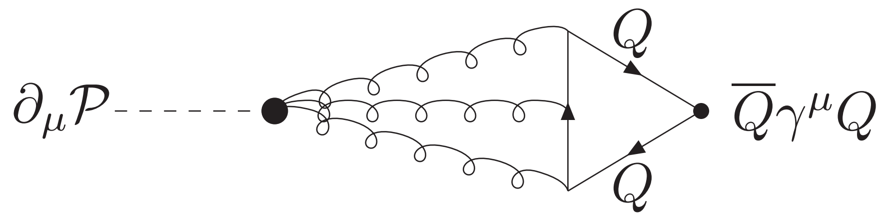


At low-energies an effective operator arises in which $\mathcal{P}$ is derivatively-coupled to three gauge bosons. This effective operator is produced via the off-diagonal coupling of $\mathcal{P}$ to the hidden sector quarks and squarks ($\mathcal{Q}$), which then rescatter and eventually (after some number of loops) convert into off-shell SM gauge bosons (gluons, photons, Ws and Zs).

The SM gauge bosons can couple to currents made up of SM fermions. This yields an effective dimension-five coupling of $\mathcal{P}$ to the SM currents, which scales as $\Lambda_{\rm ISS}^{-1}$. The dominant contribution to this effective coupling arises when the momenta of the off-shell gauge bosons are of $\mathcal{O}(\Lambda_{\rm ISS})$.

In contrast, the decay rate to on-shell gauge bosons (e.g. $\mathcal{P} \to \gamma\gamma\gamma$) is extremely suppressed due to the $\Lambda_{\rm ISS}^{-6}$ suppression of the Dolgov operator.

PNGB couplings to hadronic vector currents



Write the flavor current as a sum of the baryon number current $J_B^\mu \equiv \sum_i \bar{Q}_i \gamma^\mu Q_i$ and

$$F^\mu \sim \bar{Q}\gamma^\mu T^a Q\,, \qquad \text{where} \;\; \mathrm{Tr}\; T^a = 0\,,$$

Consider first F^μ given above. The dominant momentum passing through this diagram is of order $\Lambda_{\rm ISS}$. Thus, we can treat the Q propagators in the mass insertion approximation. An even number of mass insertions is required. The triangle with no mass insertions vanishes, and with two mass insertions yields a result proportional to $\mathrm{Tr}(M^2 T^a)$. The latter vanishes for degenerate quark masses. Hence, the above diagram yields a local interaction of the form:

$$\mathcal{L}_{\rm int} \sim \alpha_s^3(\Lambda_{\rm ISS}) \frac{\Delta m_Q^2}{\Lambda_{\rm ISS}^2} \frac{F^\mu}{\Lambda_{\rm ISS}} \partial_\mu \mathcal{P}\,.$$

Integrating by parts to get the non-derivative coupling of $\mathcal{P}$ to F^μ yields[†]

$$\mathcal{L}_{\rm int} \sim \alpha_s^3(\Lambda_{\rm ISS})\frac{\Delta m_Q^2}{\Lambda_{\rm ISS}^2}\,\mathcal{P}\,\frac{\partial_\mu F^\mu}{\Lambda_{\rm ISS}}\,.$$

At this point, we can freely shift $F^\mu \to F^\mu + cJ_B^\mu$, since $\partial_\mu J_B^\mu = 0$. As an example. consider the strangeness current $J_S^\mu = \bar{s}\gamma^\mu s$. At low momenta (much smaller than m_W),

$$\partial_\mu J_S^\mu = \frac{G_F \sin\theta_c}{\sqrt{2}}\left[\bar{u}\gamma_\nu(1-\gamma_5)d\,\bar{s}\gamma^\mu(1-\gamma_5)u + {\rm h.c.}\right]$$

due to the effective $\Delta S = \pm 1$ four-Fermi weak interaction. To compute the decay rate for $K^\pm \to \pi^\pm \mathcal{P}$, we use $\langle 0|\bar{u}\gamma_\nu(1-\gamma_5)d|\pi^-\rangle = i f_\pi q_\pi^\mu$ and similarly for K to obtain:

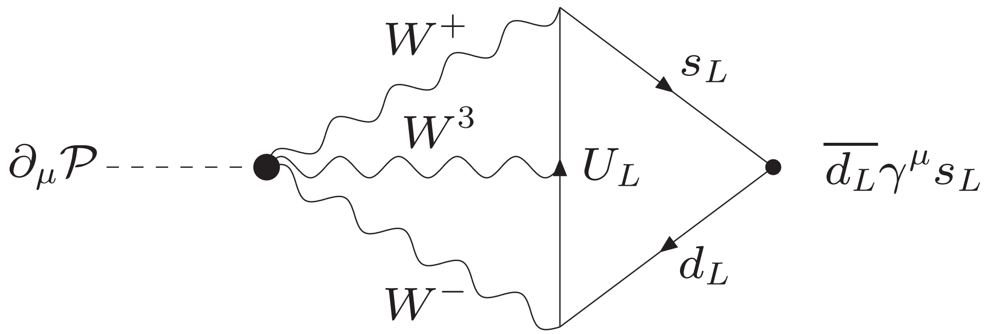
$$\mathcal{M}(K^\pm \to \pi^\pm \mathcal{P}) = \frac{G_F}{\sqrt{2}}\frac{\alpha_s^3(\Lambda_{\rm ISS})}{\Lambda_{\rm ISS}^3} m_s^2 \sin\theta_c f_\pi f_K \frac{m_K^2}{2}\,.$$

There are other contributions to the decay rate for $K^\pm \to \pi^\pm \mathcal{P}$ that are significantly larger than the one computed above. These arise from effective flavor-changing Yukawa couplings that are generated by purely weak interaction effects.

[†]If F^μ is the ordinary baryon current, then no quark mass insertions are necessary. However, after integrating by parts, we find that the low-energy coupling of $\mathcal{P}$ to the baryon current due to QCD interactions vanishes due to the conservation of the baryon current.

Flavor-violating PNGB Yukawa couplings

The flavor-violating Yukawa interactions (mediated by the electroweak sector) arise from diagrams such as:



where $U_L = (u_L, c_L, t_L)$. These contributions are GIM-suppressed by a factor of $\Delta m_q^2/\Lambda_{\rm ISS}^2$. The contribution from the top quark in the loop dominates, and the resulting effective operator for the $ds\mathcal{P}$ interaction is given by:

$$\frac{\alpha_2^3}{\Lambda_{\rm ISS}^3} m_t^2 V_{td} V_{ts}^* \overline{d}\gamma^\mu (1-\gamma_5) s\, \partial_\mu \mathcal{P} + \text{h.c.}$$

Integrating by parts yields the desired Yukawa coupling:

$$\mathcal{L}_{ds\mathcal{P}} \sim \frac{\alpha_3^2 m_t^2 V_{td} V_{ts}^* m_s}{\Lambda_{\rm ISS}^3}\, i\overline{d}\gamma_5 s\, \mathcal{P} + \text{h.c.}$$

The effective Yukawa coupling is

$$\lambda \sim \frac{\alpha_2^3 m_t^2 V_{td} V_{ts}^* m_s}{\Lambda_{\rm ISS}^3} \sim 5 \times 10^{-15} \left(\frac{2\ \rm TeV}{\Lambda_{\rm ISS}}\right)^3 ,$$

using $|V_{td}V_{ts}| \sim 3 \times 10^{-4}$. Thus,

$$\Gamma(s \to d\mathcal{P}) = \frac{\lambda^2 m_s}{16\pi} \sim 7 \times 10^{-32}\ \rm GeV \left(\frac{2\ \rm TeV}{\Lambda_{\rm ISS}}\right)^6 .$$

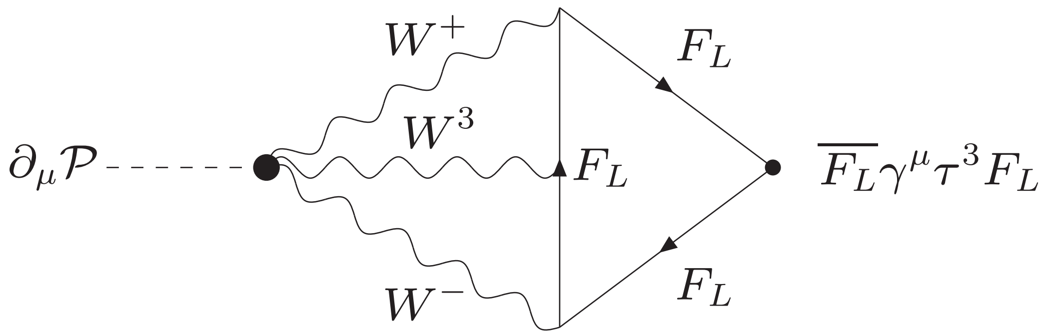
Roughly, we expect $\Gamma(K^\pm \to \pi^\pm \mathcal{P}) \sim \xi \Gamma(s \to d\mathcal{P})$, where $\xi \sim \mathcal{O}(0.1)$ due to the wave functions of the exclusive initial and final states. The $K^\pm$ lifetime is of order 10^{-8} sec., corresponding to a width of about 5×10^{-16} GeV. Thus, the branching ratio for the rare kaon decay into pion plus PNGB is roughly

$$\mathrm{BR}(K^\pm \to \pi^\pm \mathcal{P}) \sim 10^{-17} \left(\frac{2\ \rm TeV}{\Lambda_{\rm ISS}}\right)^6 .$$

The strongest experimental bounds on such decays are for approximately massless PNGBs, and give a branching ratio below 6×10^{-11}. Our predicted branching ratio is significantly below these experimental bounds.

Flavor-conserving PNGB Yukawa couplings

The flavor-conserving Yukawa interactions (mediated by the electroweak sector) arise from diagrams such as:



The resulting effective operator is of the form:

$$\frac{\alpha_2^3(\Lambda_{\rm ISS})}{\Lambda_{\rm ISS}}\overline{f}\gamma^\mu(1-\gamma_5)f\partial_\mu\mathcal{P}\,.$$

Integrating by parts and using the free field equations then yields:

$$\boxed{\mathcal{L}_{\mathcal{P}\overline{f}f} \sim \frac{\alpha_2^3(\Lambda_{\rm ISS})m_f}{\Lambda_{\rm ISS}}\, i\overline{f}\gamma_5 f\,\mathcal{P}\,,}$$

which is an allowed interaction, as the weak interactions violate C. Note that in contrast to the flavor-changing couplings that are suppressed by $1/\Lambda_{\rm ISS}^3$, the flavor-conserving couplings of $\mathcal{P}$ scale as one inverse power of $\Lambda_{\rm ISS}$.

Constraints from Astrophysics

1. The neutrino mass is nonzero, so it too will have a Yukawa coupling to $\mathcal{P}$ given by

$$\lambda_{\nu\nu\mathcal{P}} \sim \frac{\alpha_2^3 m_\nu}{\Lambda_{\rm ISS}} \sim 10^{-19}\left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right)\ .$$

This leads to a new energy loss mechanism for supernovae. Neutrinos trapped in the hot plasma can bremsstrahlung the *very* weakly interacting PNGBs, which transport energy out of the star. However, the coupling is too small for this to be a significant effect.

2. The electron Yukawa coupling $\lambda_{ee\mathcal{P}}$ can be similarly estimated. We write this in the form $\alpha_{\mathcal{P}} \equiv \lambda_{ee\mathcal{P}}^2/4\pi$:

$$\alpha_{\mathcal{P}} \sim 10^{-23}\left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right)\ .$$

The actual observational bound of Raffelt and Weiss for the coupling of a light spin zero boson (e.g., the axion) to electrons is $\alpha_a < 0.5 \times 10^{-26}$, assuming that the boson is light enough to be produced in the star (by Compton scattering or by bremsstrahlung), and assuming that it subsequently escapes. This constraint would rule out models of the type considered in this talk if $m_{\mathcal{P}} \lesssim 10^4$—$10^5$ eV and $\Lambda_{\rm ISS} \lesssim 4000$ TeV.

Conclusions

- In models of metastable SUSY-violation, one typically finds a spontaneously-broken meta-baryon number symmetry (with small explicit-breaking effects originating from very high scales). This leads to a light pseudo-Nambu-Goldstone boson, $\mathcal{P}$.

- $\mathcal{P}$ is a CP-odd, C-odd scalar. It can communicate with the Standard Model (SM) via its couplings to SM gauge bosons. C-invariance in the hidden sector dictates that the effective operator that couples $\mathcal{P}$ to SM particles must involve at least 3 intermediary SM gauge bosons.

- The only significant phenomenological constraint that we can find are from astrophysical limits on the energy loss mechanism of red giants. If the ISS scale is below 4000 TeV, this constraint imposes a $lower$ bound on the PNGB mass, which constrains the fundamental SUSY-breaking model.