The Pseudo-Nambu Goldstone Boson of Metastable SUSY-Violation

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Outline

- Directions in supersymmetry (SUSY)-breaking
- Some generic features of metastable SUSY-breaking à la ISS
- The Pseudo-Nambu-Goldstone boson (PNGB) of metastable SUSY-breaking
- C and P quantum number of the PNGB
- Spontaneous C and P violation in the metastable vacuum
- PNGB Yukawa couplings to Standard Model fermions
- Phenomenological Constraints
- Conclusions

The quest for supersymmetry breaking

- The paradigm: communication of SUSY-breaking to the MSSM

Consistent spontaneous SUSY breaking with the MSSM (or its extension) is not phenomenologically viable. So, one postulates that SUSY-breaking occurs in some other sector that is separate from the MSSM. The SUSY-breaking is then communicated to the MSSM via loops that consist of heavy particles and their superpartners of the SUSY-breaking sector.

- soft-SUSY-breaking

Integrating out the heavy particles yields the MSSM along with the requisite soft-SUSY breaking terms. The goal is to relate the soft-SUSY-breaking parameters of the MSSM to the fundamental source of SUSY-breaking.

• Viable fundamental theories of SUSY-breaking

A viable fundamental theory of SUSY-breaking (which determiness 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.

• Metastable SUSY-breaking is inevitable

To quote from ISS, metastable SUSY-breaking is inevitable with mild assumptions.

The framework for metastable SUSY-breaking

We shall employ an ISS-type model to provide the fundamental source of SUSY-breaking (in the metastable vacuum) that is communicated to 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.

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. Since these quarks carry *meta-baryon number* [we use "meta" here to distinguish this from ordinary baryon number carried by the SM quarks], the meta-baryon number 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 that violates meta-baryon number is

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

where we assume that $N_c > 3$ (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,

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

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

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

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

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

A range of possible PNGB masses

The choice of $M_{\rm ISS}$ and M_U is highly model-dependent. Since $\Lambda_{\rm ISS}$ plays the role of a messenger mass of GMSB, 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 probably as optimistic as one can get.

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

- the reduced Planck scale (2×10^{18} GeV)
- the grand unification scale (2×10^{16} GeV)
- the right-handed neutrino (seesaw) scale (10^{14} GeV)
- the axion scale (10^8—10^{12} GeV)

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

$$m_{\mathcal{P}} \sim 10^{-10} \text{ eV} \;—\; 10 \text{ GeV} \,.$$

PNGB quantum numbers

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

The couplings of $\mathcal{P}$ to the ISS sector naively conserves C and P, and 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.

Spontaneous violation of C and P

Consider the C-conserving superpotential of a toy model:

$$W = \phi_i M_{ij} \overline{\phi}_j - f^2 (M_{11} + M_{22}) \,, \qquad i, j = 1, 2 \,,$$

where the transformation laws of the superfields under C are given by

$$\phi_i \leftrightarrow \overline{\phi}_i \,, \qquad M_{ij} \leftrightarrow M_{ji} \,.$$

The scalar potential has a minimum at:

$$\langle \phi_1 \rangle = v \,, \qquad \langle \overline{\phi}_1 \rangle = \overline{v} \,, \qquad \langle M_{22} \rangle = m_0 \,,$$

with $v\overline{v} = f^2$ and all other scalar field vevs equal to zero. At the potential minimum, $V = |f^4|$ and SUSY is broken. For simplicity, we set $m_0 = 0$. We parameterize:

$$\phi_1 = v e^{i\mathcal{P}/v_0} \,, \qquad \overline{\phi}_1 = \overline{v} e^{-i\mathcal{P}/v_0} \,, \qquad v_0 \equiv \sqrt{v^2 + \overline{v}^2} \,,$$

where $\mathcal{P}$ is the Goldstone boson associated with the spontaneous breaking of meta-baryon number. Note that if $v = \overline{v}$, then C is preserved and $\mathcal{P}$ is indeed C-odd.

The fermion mass matrix and Yukawa couplings of the fermions to $\mathcal{P}$ are given by

$$\mathcal{L}_{\mathcal{P}} = \psi_{11}(v\psi_{1b} + \overline{v}\psi_1) + v\psi_{12}\psi_{2b} + \overline{v}\psi_{21}\psi_2$$
$$+ i\psi_{11}(v\psi_{1b} - \overline{v}\psi_1)\mathcal{P}/v_0 + i(v\psi_{12}\psi_{2b} - \overline{v}\psi_{21}\psi_2)\mathcal{P}/v_0 + \mathrm{h.c.}\,,$$

where $\psi_{ij} \equiv \psi_{M_{ij}}$, $\psi_i \equiv \psi_{\phi_i}$, and $\psi_{ib} \equiv \psi_{\overline{\phi}_i}$. If $v = \overline{v}$, one easily verifies that the fermion mass matrix and interactions are C-conserving. One can choose fermion mass eigenstates that are eigenstates of C. In four-component spinor notation, one easily verifies that the fermion spectrum consists of three Dirac fermions [charged with respect to an unbroken U(1)] and two neutral Majorana fermions, and $\mathcal{P}$ couples off-diagonally to fermion–antifermion pairs.

If $v \neq \overline{v}$, then C is spontaneously broken in the metastable vacuum. The three Dirac fermions are no longer eigenstates of C:

$$\Psi_{11} = \begin{pmatrix} \psi_{11} \\ \overline{\chi}_+ \end{pmatrix}, \qquad \Psi_{12} = \begin{pmatrix} \psi_{12} \\ \overline{\psi}_{2b} \end{pmatrix}, \qquad \Psi_{21} = \begin{pmatrix} \psi_{21} \\ \overline{\psi}_2 \end{pmatrix},$$

where $\chi_\pm \equiv (v\psi_{1b} \pm \overline{v}\psi_1)/v_0$.

The corresponding $\mathcal{P}$ interactions are:

$$\mathcal{L}_{\mathcal{P}\overline{\Psi}\Psi} = \mathcal{P}\left[\tfrac{1}{2}iv_0\overline{\Psi}_-(1-\gamma_5)\Psi_{11} + \text{h.c.}\right] - i\mathcal{P}\left[v\overline{\Psi}_{12}\gamma_5\Psi_{12} - \overline{v}\overline{\Psi}_{21}\gamma_5\Psi_{21}\right] ,$$

where Ψ_- is the massless four-component Majorana fermion corresponding to χ_-. The first term above exhibits the fact that P is also spontaneously violated. The second term shows the existence of diagonal pseudoscalar couplings of $\mathcal{P}$ to (a subset of the) fermion pairs. That is, $\mathcal{P}$ is now a linear combination of $J^{PC} = 0^{+-} \oplus 0^{-+}$.

One can exhibit the shift symmetry of the PNGB field more explicitly by redefining:

$$\psi_i = \tilde{\psi}_i e^{-i\mathcal{P}/v_0} , \qquad \psi_{ib} = \tilde{\psi}_{ib} e^{i\mathcal{P}/v_0} ,$$

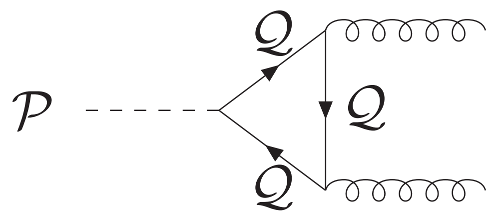
which absorbs all factors of $\mathcal{P}$ from the mass and PNGB interaction terms. Of course, the PNGB terms reappear in the kinetic energy terms, e.g.:

$$\psi_2^\dagger\overline{\sigma}^\mu\psi_2 = \frac{i}{v_0}\tilde{\psi}_2^\dagger\overline{\sigma}^\mu\tilde{\psi}_2\partial_\mu\mathcal{P} .$$

We can recover the previous non-derivative forms by integrating by parts and employing the free field equations. Of course, the results derived above also determine whether $\partial_\mu\mathcal{P}$ can couple diagonally to fermion pairs.

Coupling the PNGB to the SM sector

Since the ISS quarks carry electroweak quantum numbers, $\mathcal{P}$ can couple to the SM sector via the generic diagram:



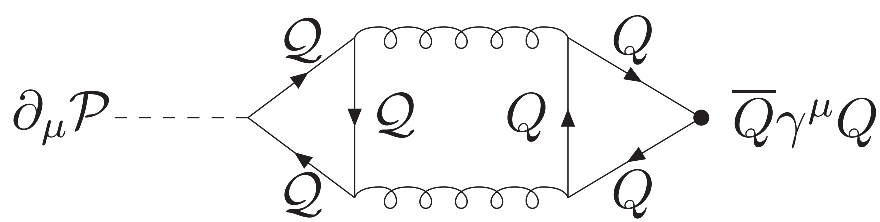
corresponding to the dimension-5 operator:

$$\mathcal{L}_{\rm eff} \sim \frac{1}{\Lambda_{\rm ISS}} \mathcal{P} F^{\mu\nu a} \widetilde{F}^a_{\mu\nu} \,,$$

as long as diagonal couplings $\mathcal{P}\mathcal{Q}\overline{\mathcal{Q}}$ are present, which requires C-breaking as exhibited above.

PNGB Yukawa couplings to SM fermions

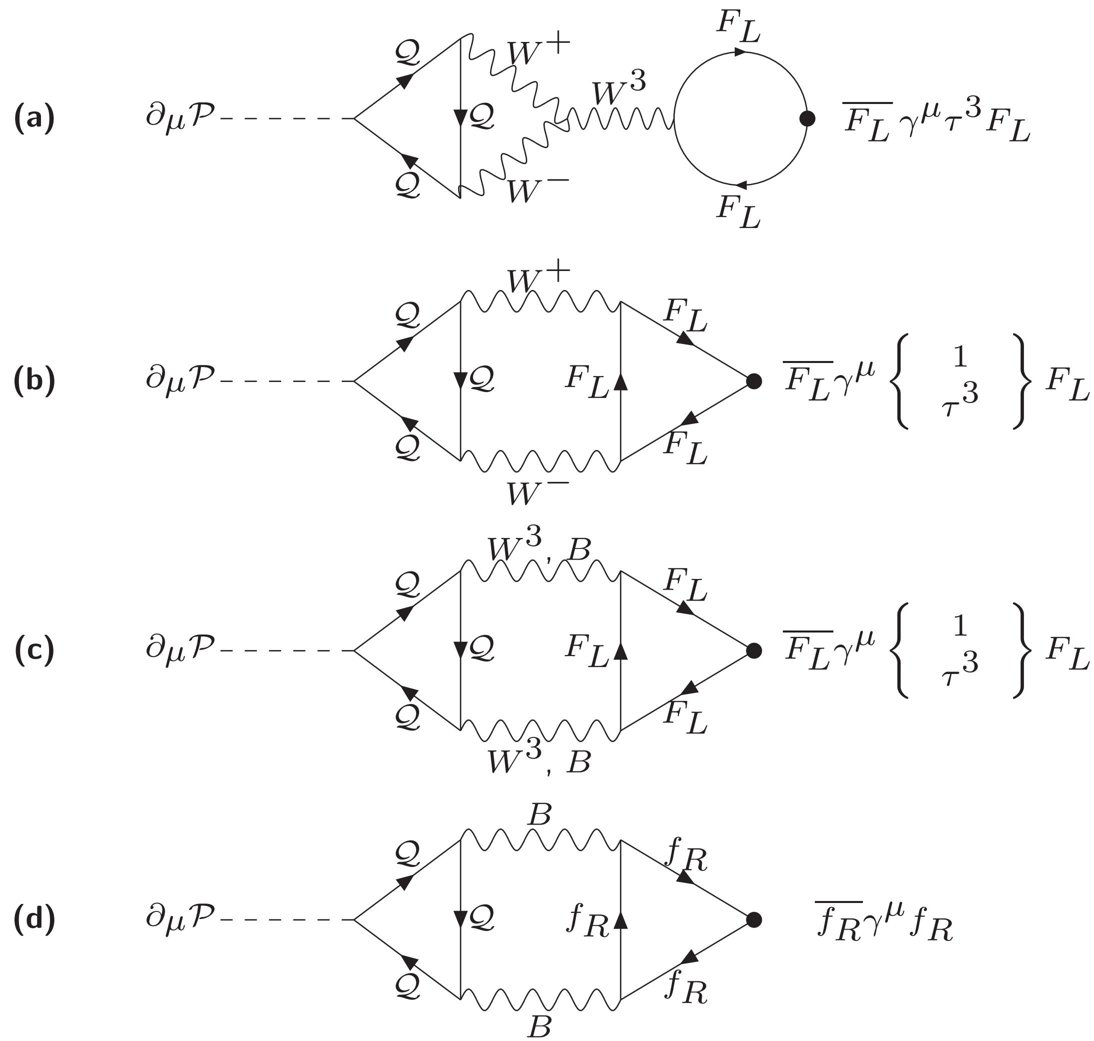
Consider first the coupling of PNGB to quarks via QCD interactions:



The diagonal couplings of $\partial_\mu \mathcal{P}$ to ISS quarks exist only if C is spontaneously broken. However, the QCD interactions preserve C, so we can use Furry's theorem to conclude that the right hand triangle vanishes (after adding up two diagrams, in which the quarks circulate in opposite directions). Three gluon exchange also vanishes, as the C properties of the diagram would imply an antisymmetric three-gluon state on the left and a symmetric three-gluon state on the right, summed over the adjoint gluon indices.

Hence, no Yukawa couplings can be generated via gluon exchange.

Next, consider PNGB Yukawa interactions mediated by the electroweak sector.



Electroweak contributions to the PNGB–fermion couplings. $F_L \equiv (u_L\,,\ d_L)$ in the case of a left-handed quark doublet, $F_L \equiv (\nu_L\,,\ e_L)$ in the case of a left-handed lepton doublet, and f_R is the right-handed quark or charged lepton singlet (with generation indices suppressed). In (c), couplings to triplet currents appear only in the mixed W^3–B box diagrams.

Again, we require spontaneous C-breaking in the ISS sector so that $\partial_\mu \mathcal{P}$ couples diagonally to the ISS quarks. However, P and C are also violated in the electroweak sector, so that all the diagram exhibited contribute.

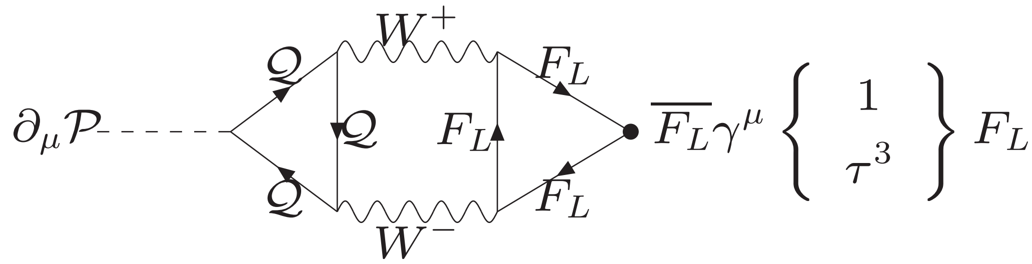
Consider first the flavor-conserving Yukawa couplings. For simplicity, set the CKM matrix to unity. Moreover, as $g_2 > g_1$, we will keep only the leading contribution that is proportional to α_2^2. The loops are dominated by the ISS scale, hence we expect an effective operator of the form:

$$\frac{\alpha_2^2(\Lambda_{\rm ISS})}{\Lambda_{\rm ISS}} \overline{f} \gamma^\mu (1-\gamma_5) f \partial_\mu \mathcal{P} \ .$$

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

$$\boxed{\mathcal{L}_{\mathcal{P}\overline{f}f} \sim \frac{\alpha_2^2(\Lambda_{\rm ISS}) m_f}{\Lambda_{\rm ISS}} \, i\overline{f}\gamma_5 f \, \mathcal{P} \ .}$$

Next, we examine the flavor-changing Yukawa coupling, which can arise from:



Inserting the factors of the CKM matrix V at the two W vertices of the right-hand triangle, we note that if the quark masses were degenerate one would produce a factor of $V^\dagger V = \mathbb{1}$. Thus, we have a GIM suppression. Treating the quark masses in the mass-insertion approximation, one needs two mass insertions to obtain a flavor-changing coupling. This yields a suppression factor of $\Delta m_q^2/\Lambda_{\rm ISS}^2$. The top quark in the loop should dominate, and the resulting effective operator for $\overline{d}s\mathcal{P}$ is:

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

Integrating by parts yields the desired Yukawa coupling:

$$\boxed{\mathcal{L}_{\overline{d}s\mathcal{P}} \sim \frac{\alpha_2^2 m_t^2 V_{td} V_{ts}^* m_s}{\Lambda_{\rm ISS}^3} (i\overline{d}\gamma_5 s\, \mathcal{P} + \mathrm{h.c.})\ .}$$

Constraints from $K^\pm \to \pi^\pm \mathcal{P}$

We compute the decay width for $s \to d\mathcal{P}$. The effective Yukawa coupling is

$$\lambda \sim \frac{\alpha_2^2 m_t^2 V_{td} V_{ts}^* m_s}{\Lambda_{\rm ISS}^3} \sim 2 \times 10^{-13} \left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right)^3 ,$$

where I have used $\alpha_2^2 \sim 10^{-3}$ and $|V_{td}V_{ts}| \sim 3 \times 10^{-4}$. Thus, for $\Lambda_{\rm ISS} = 2$ TeV,

$$\Gamma(s \to d\mathcal{P}) = \frac{\lambda^2 m_s}{64\pi^2} \sim 7 \times 10^{-30}\ {\rm GeV} .$$

If we crudely use this result for $\Gamma(K^\pm \to \pi^\pm \mathcal{P})$, then we would obtain the predicted branching ratio by dividing out by the total decay width, $\Gamma(K) \sim 5 \times 10^{-16}$ GeV. Thus, we would expect:

$$\mathrm{BR}(K^\pm \to \pi^\pm \mathcal{P}) \sim 10^{-14} \left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right)^6 ,$$

which should be compared with the current experimental bound of 6×10^{-11}.

Constraints from $\Upsilon \to \gamma\mathcal{P}$

If V is a ${}^3S_1$ quarkonium bound state of $\overline{Q}Q$, and if the $\mathcal{P}\overline{Q}Q$ coupling is given by $\lambda_Q m_Q i\overline{Q}\gamma_5 Q\mathcal{P}$, then a tree-level computation yields:

$$\frac{\Gamma(V \to \gamma\mathcal{P})}{\Gamma(V \to e^+e^-)} = \frac{\lambda_Q^2 m_V^2}{8\pi\alpha}\,.$$

We have estimated $\lambda_Q \sim \alpha_2^2/\Lambda_{\rm ISS}$. Hence, using BR$(\Upsilon \to e^+e^-) = 2\%$, we find:

$$BR(\Upsilon \to \gamma\mathcal{P}) \sim 2 \times 10^{-12}\left(\frac{2\ \mathrm{TeV}}{\Lambda_{\rm ISS}}\right)^2\,.$$

which is many orders away from experimental bounds. The corresponding branching ratio for $\psi \to \gamma\mathcal{P}$ is about a factor of three smaller.

Constraints from Astrophysics

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

$$\lambda_{\nu\nu\mathcal{P}} \sim \frac{\alpha_2^2 m_\nu}{\Lambda_{\rm ISS}} \sim 10^{-17}\left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right) .$$

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 $\alpha_{\mathcal{P}} \equiv \lambda_{ee\mathcal{P}}^2/4\pi$:

$$\alpha_{\mathcal{P}} \sim 10^{-20}\left(\frac{2\ {\rm TeV}}{\Lambda_{\rm ISS}}\right) .$$

The actual observational bound of Raffelt and Weiss for the coupling of a light spin zero boson 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 that is 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 .

Conclusions

- Metastable SUSY-violation may be a critical ingredient for constructing theories of fundamental SUSY-breaking. In ISS-type models, one typically finds a spontaneously-broken meta-baryon number symmetry (with small explicit-breaking effects from very high scales). This leads to a light pseudo-Nambu-Goldstone boson, $\mathcal{P}$.

- $\mathcal{P}$ is a CP-odd, C-odd scalar. If C and P are spontaneously broken in the metastable vacuum, then it is possible for $\mathcal{P}$ to couple diagonally to ISS-quarks. Assuming that these quarks also possess SM gauge quantum numbers, interactions of $\mathcal{P}$ with SM particles are generated.

- The only significant phenomenological constraint that we can find are from astrophysical limits on the energy loss mechanism of red giants. This constraint imposes a $lower$ bound on the PNGB mass, which places restrictions on the parameters of the SUSY-breaking model.