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Phenomenology of Light Gauginos
I. Motivation, Masses, Lifetimes and Limits

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Abstract: I explore an economical variant on supersymmetric standard models which may be indicated on cosmological grounds and is shown to have no SUSY-CP problem. Demanding radiative electroweak symmetry breaking suggests that the Higgs is light; other scalar masses may be $\sim 100-200$ GeV or less. In this case the gluino and photino, while massless at tree level, have 1-loop masses $m_{\tilde{g}} \sim 100-600$ MeV and $m_{\tilde{\gamma}} \sim 100-1000$ MeV. New hadrons with mass $\sim 1-3$ GeV are predicted and their lifetimes estimated. Existing experimental limits are discussed.

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The customary approach to studying the phenomenological implications of supersymmetry has been to assume that the "low energy" effective Lagrangian contains all possible renormalizable operators, including in principle all possible soft supersymmetry breaking terms, consistent with the gauge symmetries and certain global and discrete symmetries. Some models of SUSY-breaking naturally lead to relations among the SUSY-breaking parameters at the scale M_{SUSY}, so that, e.g., the minimal supersymmetric standard model (MSSM) requires specification of 6-8 parameters beyond the gauge and Yukawa couplings already determined in the MSM: $tan\beta \equiv \frac{v_U}{v_D}$, the ratio of the two Higgs vevs; μ, the coefficient of the SUSY-invariant coupling between higgsinos; M_0, a universal SUSY-breaking scalar mass; m_{12}^2, the SUSY-breaking mixing in the mass-squared matrix of the Higgs scalars (μB or $\mu M_0 B$ in alternate notations); $M_{1,2,3}$, the SUSY-breaking gaugino masses (proportional to one another if the MSSM is embedded in a GUT); and A, the coefficient of SUSY-breaking scalar trilinear terms obtained by replacing the fermions in the MSM Yukawa terms by their superpartners. To obtain predictions for the actual superparticle spectrum in terms of these basic parameters, the renormalization group equations for masses, mixings and couplings are evolved from the scale M_{SUSY} to the scale m_{Z^0} where on account of different RG running and flavor dependent couplings, the various scalars and fermions can have quite different masses. A particularly attractive aspect of this approach is that for a heavy top quark, the mass-squared of the Higgs field which gives mass to the charge 2/3 quarks becomes negative at low energy and the electroweak symmetry is spontaneously broken[1, 2], with m_{Z^0} a function of A, M_0 and other parameters of the theory. In this conventional treatment of the MSSM, the lightest squark mass is constrained by experiment to be greater than 126 GeV and the gluino mass to be greater than 141 GeV, given favorable assumptions regarding their decays[3].

I will argue here that a more restrictive form of low energy SUSY breaking may actually be used by Nature, one without dimension-3 operators, i.e., with no tree level gaugino masses or scalar trilinear couplings. We shall see that the remaining parameters of the theory are rather well-constrained when electroweak symmetry breaking is demanded, and that the resultant model is both very predictive and consistent with laboratory and cosmological observations. If this is the correct structure of the low energy world, there will be many consequences which can be discovered and investigated before the construction of the LHC. Some of these are discussed in a companion paper[4][2], hereafter refered to as II. The purpose of this Letter is threefold:

1. Articulate the theoretical motivation for the absence of D=3 SUSY breaking operators in the MSSM.

2. Focus on the most probable portion of (M_0, μ, $tan\beta$) parameter space to obtain predictions for the gluino and photino masses.

3. Determine the mass and lifetime of the lightest R-hadrons, and with that information establish the experimental limits relevant to this scenario.

There are several reasons to suspect that there may be no dimension-3 SUSY breaking operators in the low energy theory. Firstly, their absence accounts for the absence of an observable neutron electric dipole moment and other CP violating effects which arise naturally with conventional SUSY breaking (the "SUSY CP problem"). With no dimension-3 SUSY breaking operators, the only CP violating phases not already present in the MSM are in the terms[3] $\mu \int \hat{H}_U \hat{H}_D d^2\theta$ and $m_{12}^2 H_U H_D$, and possibly in the SUSY-breaking scalar mass-squared matrices if they are flavor non-diagonal. Using an R-transformation and a $U(1)_{PQ}$ phase rotation on the superfield $(\hat{H}_U \hat{H}_D)$, phases in both μ and m_{12}^2 can be removed. Any phase which is introduced thereby into the Yukawa terms in the superpotential can be removed by chiral

[2]A preliminary discussion of many points developed here and in (II) was given in [5].
[3]Hatted fields denote superfields, unhatted fields their scalar components.

transformations on the quark superfields, merely changing the phases which contribute to the strong CP problem (which must be solved by some other mechanism). Since the gauge-kinetic terms are not affected by U(1) and R transformations, the preceding manipulations do not introduce phases in interactions involving gauginos. Hence we see that without dimension-3 SUSY breaking, the only phases beyond those of the MSM are in the squark mass-squared matrices. However to generate an electric dipole moment, which is a chirality flip operator, requires a phase in an off-diagonal term mixing superpartners of left and right chiral quarks. In the case at hand, mixing between left and right chiral superpartners is induced only by the μ term and not also by A terms as in the usual case. Therefore the relevant mixing is real in the basis found above. Note that the argument given here applies to all orders of perturbation theory, so it shows that even though the physical gluino and neutralinos do have a mass coming from radiative corrections, no edm is generated unless A terms are present.[4] Finding a neutron edm would therefore exclude the scenario of no dimension-3 SUSY breaking.

A second reason to consider the absence of dimension-3 SUSY breaking terms is that they simply do not arise in many types of SUSY breaking. The reason a distinction naturally emerges between dimension-3 and dimension-2 SUSY breaking can be seen as follows. A SUSY-breaking mass for the spin-0 component of a chiral superfield $\hat{Q}$ originates from its kinetic term:

$$\int K(\hat{\Psi}_i{}^\dagger, \hat{\Psi}_i)\hat{Q}^\dagger\hat{Q} d^2\theta d^2\bar{\theta}, \tag{1}$$

where the Ψ_i includes all the chiral superfields of the theory. The Kahler potential $K(\hat{\Psi}_i{}^\dagger\hat{\Psi}_i)$ is a vector superfield, so generally has an expansion $1 + \frac{b_i}{M}(\hat{\Psi}_i + \hat{\Psi}_i{}^\dagger) + \frac{c_i}{M^2}\hat{\Psi}_i{}^\dagger\hat{\Psi}_i + \ldots$. In a hidden sector SUSY-breaking scenario the interaction between hidden sector and visible sector fields is purely gravi-

[4] A discussion of natural criteria for eliminating the SUSY CP problem in the MSSM, including the case that scalar trilinears and gaugino masses are present, can be found in [6].

tational so that $M \sim M_{Pl}$ for terms coupling visible and hidden sector fields. When one or more of the hidden sector superfields develops a non-vanishing auxilliary component, a mass-squared $M_0^2 \sim c_i(\frac{F_{\Psi_i}}{M_{Pl}})^2$ for the scalar component of $\hat{Q}$ is produced. On the other hand, gaugino masses come from the superpotential:

$$\int f(\hat{\Psi}_i) W_\alpha W^\alpha d^2\theta, \tag{2}$$

where the gauge kinetic function $f(\hat{\Psi}_i)$ is a gauge singlet chiral superfield whose expansion in hidden sector fields has the form $1+\frac{b_i}{M_{Pl}}\hat{\Psi}_i+\frac{c_{ij}}{M_{Pl}^2}\hat{\Psi}_i\hat{\Psi}_j+\ldots$. If the linear terms in this expansion have no $F-$component or are absent entirely, for example because there are no gauge singlet hidden sector fields, then the leading contribution to the gaugino mass is $\sim \frac{c_{ij}F_{\Psi_i}<\Psi_j>}{M_{Pl}^2} \sim \frac{<\Psi_j>}{M_{Pl}}M_0$. As discussed in detail in ref. [7], hidden sector models only make sense if $< \Psi' > \;\; << \;\; M_{Pl}$, so the dimension-3 gaugino mass is negligible compared to the dimension-2 scalar masses. A terms are produced in the same way as the gaugino mass, replacing $W_\alpha W^\alpha$ in eqn (2) by the Yukawa terms, or by linear terms in the Kahler potential after using the equation of motion to eliminate the F-component of $\hat{Q}$ or $\hat{Q}^\dagger$.

Thus in hidden sector SUSY-breaking models in which gauge singlets do not develop an F-component, the coefficients of dimension-3 operators are negligible in comparison to the coefficients of dimension-2 operators. More generally, this occurs whenever linear terms which develop a non-vanishing F-component are absent from the expansion of the Kahler potential, gauge kinetic function and the analogous functions for other terms in the superpotential. This occurs in several models, for instance ones in which the cosmological constant naturally vanishes in leading order[8] and others in which SUSY-breaking is driven by hidden sector gaugino condensation and the effective Lagrangian is invariant under a phase transformation on the condensate.

The success of standard cosmology and nucleosynthesis may be another hint that SUSY-breaking is not driven by gauge singlet fields since such fields generally produce severe cosmological difficulties as shown in ref. [7][5].

Anticipating results to be obtained below, a final motivation for considering this scenario is that it gives a natural explanation for the missing matter of the Universe. For R-hadron and photino masses in the ranges predicted[5] in this scenario, relic photinos provide the observed level of cold dark matter in the Universe[10]. In particular pions catalyze the conversion of photinos to R^0's, the gluon-gluino bound state, which annihilate via the strong interactions. For a critical value of $r \equiv \frac{m(R^0)}{m_{\tilde{\gamma}}}$ in the range $\sim 1.6 - 2.2$, the resultant density of photinos is just what is needed.

Therefore we henceforth make the ansatz that there are no dimension-3 SUSY-breaking operators, and set all A's and $M_{1,2,3}$ to zero. The gluino and lightest neutralino, which are massless in tree approximation, get masses at one loop from virtual top-stop pairs, and, for the neutralinos, from "electroweak" loops involving higgsino/wino/bino and Higgs/gauge bosons[11, 12, 13]. The top-stop loop depends on the stop masses, especially the splitting between the stop mass eigenstates, which is proportional to $\mu cot\beta$ and the average stop mass. The electroweak loops depend on the Higgs and Higgsino masses and mixings, especially on μ, $\tan\beta$, and the masses of the heavier Higgs bosons. These radiative corrections were estimated in ref. [13], in the limit of $\mu,\ M_0 >> M_Z$, assuming a common scalar mass and taking various values of M_0, μ, and $tan\beta$. There, it was determined that in order to insure that the chargino mass is greater than its LEP lower bound of about 45 GeV, μ must either be less than 100 GeV (and $tan\beta \lesssim 2$) or greater than several TeV.

Here I will also suppose that radiative electroweak symmetry breaking[1, 2] produces the observed Z^0 mass for $m_t \sim 175$ GeV. This is not possible in the large μ region, so I will consider only the low μ region: $\mu \lesssim 100$ GeV. In

[5]In special situations the difficulties can be overcome, as shown in ref. [9].

addition, from Fig. 6 of ref. [1] one sees that M_0, the SUSY-breaking scalar mass in the Higgs sector must be $\sim 100 - 300$ GeV, with 150 GeV being the favored value. A more exact treatment suggests a somewhat lower value. Assuming that the stop mass is similar to M_0, from Figs. 4 and 5 of ref. [13] we find $m_{\tilde{g}} \sim 100 - 600$ and $m_{\tilde{\gamma}} \sim 100 - 900$ MeV.[6] Since the electroweak loop was treated in ref. [13] with an approximation which is valid when M_0 and μ are $>> m_{Z^0}$, these results for the photino mass are only indicative of the range to be expected. Furthermore, in order to properly take into account the differences between masses of various squarks and the parameter M_0, a more detailed treatment is required. For the present, I will attach a $\sim$factor-of-two uncertainty to the electroweak loop and consider the enlarged photino mass range $100 - 1000$ MeV.

Having outlined above the motivation for considering theories without dimension-3 SUSY breaking operators and having focused on a substantially restricted range of parameters, let us turn to consideration of the most essential phenomenological properties of the light particles of this theory. The primary issues to be here discussed are: *i)* Predicted mass and lifetime of the lightest R-meson, the $g\tilde{g}$ (glueballino) bound state denoted R^0. *ii)* Predicted mass of the flavor singlet pseudoscalar which gets its mass via the anomaly (the "extra" pseudoscalar corresponding to the $\tilde{g}\tilde{g}$ ground state degree of freedom). *iii)* The flavor singlet pseudogoldstone boson resulting from the spontaneous breaking of the extra chiral symmetry associated with the light gluino, which is identified as the η'.

The R^0 mass can be quite well determined from existing lattice QCD calculations, as follows[16]. If the gluino were massless and there were no

[6] Imposing strict equality of all scalar SUSY-breaking masses at the scale M_{SUSY} is difficult or maybe impossible, since in that case the lightest Higgs mass comes out too low given the restrictions on μ, $tan\beta$ coming from chargino and neutralino masses[14, 15]. However radiative corrections are sufficient to give an acceptable mass to the lightest Higgs if the stop mass is allowed to be larger than this. This favors the low end of the gluino mass range while not much affecting the photino mass prediction.

quarks in the theory (let us call this theory sYM for super-Yang Mills), SUSY would be unbroken and the R^0 would be in a degenerate supermultiplet with the lightest 0^{++} glueball, G, and a 0^{-+} state I shall denote $\tilde{\eta}$, which can be thought of as a $\tilde{g}\tilde{g}$ bound state[7]. To the extent that quenched approximation is accurate for sYM,[8] the mass of the physical R^0 in the continuum limit of this theory would be the same as the mass of the 0^{++} glueball. The latest quenched lattice QCD value of $m(G)$ from the GF11 group is 1740 ± 71 MeV[17]. Note the increase from the 1440 ± 110 value given in ref. [18] and used in my earlier work[16, 5]. The UKQCD collaboration reports[19] 1550 ± 50 MeV for the 0^{++} mass, but this error is only statistical. Adding in quadrature a 70 MeV lattice error and a 15% quenching uncertainty[9] leads to a total uncertainty of $\sim$ 270 MeV, so I will use the range 1.3 - 2 GeV for massless gluinos. Experimentally, the $f_0(1520)$ and $f_0(1720)$ seem to be the leading candidates for the ground state glueball, but the situation is still unclear.

Physical flavor singlet 0^{++} states in the glueball mass region will in general contain both glueball and $q\bar{q}$ components, causing physical masses not to correspond to the lattice value. While mixing with other states causes the physical glueball and the $\tilde{\eta}$ to shift, the R^0 has nothing nearby with which to mix. Thus the sYM glueball mass may give a better estimate of the R^0 mass than it does of the physical 0^{++} masses. In analogy with the dependence of baryon mass on quark mass, we can expect $m(R^0) \approx m(G) + m_{\tilde{g}}$, where

[7]It is convenient to think of the states in terms of their "valence" constituents but of course each carries a "sea" so, e.g., the glueball may be better described as a coherent state of many soft gluons than as a state of two gluons. Knowledge of these aspects of the states is not needed for estimation of their masses.

[8]The 1-loop beta function is the same for sYM as for ordinary QCD with 3 light quarks, so the accuracy estimate for quenched approximation in ordinary QCD, $5 - 15\%$, should be applicable here.

[9]The uncertainty associated with quenched approximation with both light quarks and gluinos was taken in [16] to be 25%. However since the estimate of the quenching error for ordinary QCD is obtained by comparing lattice results with the hadron spectrum, it already includes the effects of gluinos, if they are present in nature.

$m(G)$ is the unmixed glueball mass. Therefore in view of the expected small gluino mass and the various uncertainties discussed above, I shall adopt the estimate $1.4-2.2$ GeV for the R^0 mass, while giving greatest credence to the range $1.5-2$ GeV.[10]

In sYM, which in quenched approximation is identical to ordinary QCD, the $\tilde{\eta}$ with mass $\sim 1\frac{1}{2}$ GeV is the pseudoscalar that gets its mass from the anomaly. Thus in QCD with light gluinos the particle which gets its mass from the anomaly is too heavy to be the η'. However there is a non-anomalous chiral U(1) formed from the usual chiral U(1) of the light quarks and the chiral R-symmetry of the gluinos[16]. Due to the formation of $q\bar{q}$ and $\tilde{g}\tilde{g}$ condensates, $< \bar{q}q >$ and $< \bar{\lambda}\lambda >$, this chiral symmetry is spontaneously broken. Therefore, it is natural to identify the η' with the pseudogoldstone boson associated with the spontaneously broken U(1). Using the usual PCAC and current algebra techniques, in ref. [16] I obtained the relationship between masses and condensates necessary to produce the correct η' mass (ignoring mixing): $m_{\tilde{g}} < \bar{\lambda}\lambda > \quad \sim 10 \ m_s < s\bar{s} >$. The required gluino condensate is reasonable, for $m_{\tilde{g}} \sim 100-300$ MeV.[11] In a more refined discussion, the physical η' would be treated as a superposition of the pseudo-goldstone boson, the orthogonal state which gets its mass from the anomaly, and the η. I have not yet identified any clear test for the prediction that the η' is a

[10] A dedicated lattice gauge theory calculation of the masses of these particles could in principle improve these estimates. Such a calculation has the usual difficulty of treating chiral fermions on the lattice, due to the Majorana nature of the gluino. On the other hand, since the R^0 does not have vacuum quantum numbers, some of the difficulties in a glueball mass calculation are absent.

[11] Note that ensuring $m_{\tilde{g}} \gtrsim 100$ MeV requires the stop squarks to be not too heavy, or else their fractional splitting is too small given that the off-diagonal term in the squark mass matrix, $\mu cot\beta$, is limited. This requires the average value of the stop mass M_{st} to be $\lesssim 300$ GeV[13], which is of the same order as the value $M_0 \sim 150$ GeV indicated by electroweak symmetry breaking. For $M_{st} = 150$ GeV, ensuring $m_{\tilde{g}} \gtrsim 100$ MeV requires $\mu \gtrsim 40$ GeV for $tan\beta = 2$ and $\mu \gtrsim 20$ GeV for $tan\beta = 1$. This eliminates the otherwise attractive strategy of requiring μ to arise from SUSY-breaking. In this scenario that would cause $\mu = 0$ due to its being a dimension-3 term. This would solve the strong CP problem but replace it with the old U(1) problem.

pseudogoldstone boson and contains a $\sim$ 30% $\tilde{g}\tilde{g}$ component, since model independent predictions concerning the η' are for ratios in which the gluino component plays no role.

An important point, independent of details of the mixing, is that this scenario *predicts* the existance of a flavor singlet pseudoscalar meson in addition to the η' which is not a part of the conventional QCD spectrum of quark mesons and glueballs, whose mass should be in the $1\frac{1}{2} - 2$ GeV range, apart from mixing. A detailed discussion of this and other flavor singlet mesons will be left for the future. Note however that the isosinglet pseudoscalar at 1420 MeV discovered by MarkIII[20] and DM2[21] in radiative J/Ψ decay and recently confirmed by the Crystal Barrel in $p\bar{p}$ annihilation[22], is incompatible with any conventional quark model (the closest quark-model multiplet with an opening has a pion mass of 1800 MeV) or glueball interpretation[19] and appears to be an excellent candidate for the expected extra state[23].

Having in hand an estimate of the R^0 mass and photino mass, we now return to determining the R^0 lifetime. Making an absolute estimate of the lifetime of a light hadron is always problematic. Although the relevant short distance operators can be accurately fixed in terms of the parameters of the Lagrangian which we have constrained to a considerable extent, hadronic matrix elements are difficult to determine. It is particularly tricky for the R^0 in this scenario because the photino mass is larger than the current gluino mass and, since $m_{\tilde{\gamma}} \sim \frac{1}{2} m_{R^0}$, the decay is highly suppressed even using a constituent mass for the gluino. The decay rate of a free gluino into a photino and massless $u\bar{u}$ and $d\bar{d}$ pairs is known[24]:

$$\Gamma_0(m_{\tilde{g}}, m_{\tilde{\gamma}}) = \frac{\alpha \alpha_s m_{\tilde{g}}^5}{48\pi M_{sq}^4} \frac{5}{9} f(\frac{m_{\tilde{\gamma}}}{m_{\tilde{g}}}), \tag{3}$$

taking M_{sq} to be a common up and down squark mass. The function $f(y) = [(1-y^2)(1+2y-7y^2+20y^3-7y^4+2y^5+y^6)+24y^3(1-y+y^2)log(y)]$ contains the phase space suppression which is important when the photino is massive. The problem is to take into account how interactions with the gluon and "sea" inside the R^0 "loans" mass to the gluino. If this effect is ignored one would find the R^0 to be absolutely stable except for the largest gluino and smallest photino masses.

A method of estimating the *maximal* effect of such a "loan", and thus a lower limit on the R^0 lifetime, can be obtained by elaborating a suggestion of refs. [25, 26]. The basic idea is to think of the hadron (here the R^0) as a bare massless parton (in this case a gluon) carrying momentum fraction x and a remainder (here, the gluino) having an effective mass $M\sqrt{1-x}$, where M is the mass of the decaying hadron. Then the structure function, giving the probability distribution of partons of fraction x, also gives the distribution of effective masses for the remainder (here, the gluino). Summing the decay rate for gluinos of effective mass $m(R^0)\sqrt{1-x}$ over the probability distribution for the gluino to have this effective mass, leads to a crude estimate or upper bound on the rate:

$$\Gamma(m(R^0), z) = \Gamma_0(m(R^0), 0) \int_0^{1-z^2} (1-x)^{\frac{5}{2}} F(x) dx f(z/\sqrt{1-x}), \qquad (4)$$

where $z = \frac{m_{\tilde{\gamma}}}{m(R^0)}$. The distribution function of the gluon in the R^0 is unknown, but can be bracketed with extreme cases: the non-relativistic $F_{nr}(x) = \delta(x - \frac{1}{2})$ and the ultrarelativistic $F_{ur}(x) = 6x(1-x)$. The normalizations are chosen so that half the R^0's momentum is carried by gluons. Figure 1 shows the R^0 lifetime produced by this model, for $M_{sq} = 150$ GeV and $m(R^0) = 1.5$ GeV, for these two structure functions, and also for the intermediate choice $F_{10}(x) = N_{10}x^{10}(1-x)^{10}$, as a function of $r \equiv z^{-1} = \frac{m(R^0)}{m_{\tilde{\gamma}}}$. Results for any R^0 and squark mass can be found from this figure using the scaling behavior $\Gamma(m(R^0), M_{sq}, z) \sim m(R^0)^5 M_{sq}^{-4} g(z)$, as long as it is legitimate to ignore the mass of the remnant hadronic system, say a pion.

The decay rates produced in this model can be considered upper limits on the actual decay rate, because the model maximizes the "loan" in dynamical mass which can be made by the gluons to the gluino. We can get an idea of the accuracy of this model by using it to estimate the kaon semileptonic decay rate. $K_{\mu 3}$ decay presents a similar dynamical problem to $R^0 \rightarrow \tilde{\gamma} + X$ since $m_\mu \sim m_s$. (The problem is more severe for $R^0 \rightarrow \tilde{\gamma} + X$ since the photino is expected to be heavier than the gluino, and also the mass ratio $m(R^0)/m_{\tilde{\gamma}}$ is probably less than the ratio m_K/m_μ.) This model gives an approximately correct ratio between $K_{\mu 3}$ and K_{e3} rates: 0.72 or 0.81 for the non-relativistic and ultra-relativistic wavefunctions, respectively, compared to the experimental value of 0.67. However it predicts a $K_{\mu 3}$ rates 2-4 times larger than observed, for the same two wavefunctions, overestimating the rate as anticipated. Since the non-relativistic wavefunction gives better predictions for both quantities, we will favor its predictions for the R^0 lifetime.

In ref. [16] I reported the result of a comprehensive study of relevant experiments, including all those used in the famous UA1 analysis[27] which has widely been accepted as excluding all but certain small "windows" for low gluino mass. As noted in [16], the UA1 analysis *assumed* that the gluino lifetime is short enough that missing energy and beam dump experiments are sensitive to it. However R-hadrons produced in the target or beam dump degrade in energy very rapidly due to their strong interaction scattering length of $\sim 10 - 15$ cm. Since the photino is supposed to reinteract in the detector downstream of the beam dump or carry off appreciable missing energy, it typically has enough energy to be recognized only if it is emitted before the R-hadron interacts. As discussed in connection with a particular experiment in ref. [28], and more generally in ref. [16], if the R^0 lifetime is longer than $\sim 5\ 10^{-11}$ sec this criterion is not met and the sensitivity of beam dump and missing energy experiments to light gluinos is degraded.

Although the R^0 lifetime estimate obtained above has a large uncertainty, for nearly all of the parameter space of interest the lower bound on the lifetime is long enough that we must deal with the degradation issue. By a mild theoretical idealization, we can treat the effect of a finite lifetime analytically. Suppose that all R^0's are produced with the same energy so that each of them has the same time dilation factor $\gamma = \frac{E}{m(R^0)}$. Then the ratio of the probability of the R^0 producing a photino before interacting, compared to what it would be if the lifetime were zero is:

$$p(\tau) = (1 + \frac{\gamma\beta c\tau}{\lambda})^{-1}, \tag{5}$$

taking the R^0 lifetime to be τ and its interaction length to be λ (which is approximately the same as for a nucleon, so $\lambda \sim 10-15$ cm). The reduction in the expected number of events when the R^0 lifetime is non-zero corresponds to a reduction in sensitivity to squark mass by a factor $p(\tau)^{-1/4}$. In the BEBC experiment[29], $< \gamma > \beta c \sim 1.2 \times 10^{12}$ cm/s. This experiment modeled the loss due to rescattering in the dump for a given $m_{\tilde{g}}$ and M_{sq}, taking τ to be the lifetime for a free gluino to decay to a massless photino and $u\bar{u}$, $d\bar{d}$ or $s\bar{s}$ pair. This is an appropriate procedure for the portions of parameter space in which the gluino is much heavier than the photino and its mass is much larger than the the confinement scale (say, $m_{\tilde{g}} \gtrsim 2$ GeV). However the photino emission time obtained in this way is much shorter than when the gluino is actually light and inside a massive R^0, and the R^0 lifetime is suppressed on account of the photino mass. From the BEBC figure, their squark mass limit is ~ 330 GeV, for a "gluino" mass (effectively, $m(R^0)$)[12] of ~ 1.7 GeV. This corresponds to a lifetime of 10^{-10} sec using their formula[13]. Therefore their squark mass limit for a lifetime 10^{-9} (10^{-8}, 10^{-7}) sec becomes 185 GeV (107, 60) GeV rather than 330 GeV. Note that this is essentially a limit on the mass of the lightest u squark because the photino couples to charge. The d squark could be a factor $\sim \frac{1}{\sqrt{2}}$ lighter.

In (II) I show that the experiment of Bernstein et al[30] is actually insensitive to an R^0 in the interesting range of masses. Combining these new

[12] I thank A. Cooper-Sarkar for correspondence on this point.

[13] Which gives a factor 1.8 larger lifetime than obtained using eq. (3) with $\alpha_s = 0.117$, since they take $\alpha_s = 0.15$ and allow decay into $s\bar{s}$ pairs which is kinematically forbidden in the parameter range of interest here.

facts with the analysis of ref. [16] (where references are given), leads to Fig. 2, showing the excluded regions for the R^0 mass-lifetime plane. ARGUS gives the light grey region, assuming $m(R^0) = 1.5$ GeV; CUSB gives the next-to-darkest block, with its excluded region extending over all lifetimes. Gustafson et al gives the next-to-lightest block in the upper portion of the figure; it extends to infinite lifetime, but makes specific assumptions about production rate. UA1 gives the darkest block in the lower right corner; it extends to higher masses and shorter lifetimes not shown on the figure, where it is continued by collider limits. Evidently, the most interesting regions for the tree-level-massless gluino scenario are essentially unconstrained by previous experiments.

The phenomenology discussed above also applies to theories with a small tree-level gluino mass. Compatibility with the η' mass and the CUSB experiment requires 100MeV $\lesssim m_{\tilde{g}} \lesssim 1.5$ GeV. The photino mass would have to be tuned to be close enough to the R^0 mass to avoid too much relic density in photinos[31]. The extent of the required tuning increases as the squark mass does. It would be very difficult to have gluinos heavy enough to avoid the CUSB limit, $m_{\tilde{g}} \gtrsim 3.5$ GeV, while keeping the lifetime short enough to avoid conflict with missing energy experiments. Thus the gluino mass must be either less than 1 GeV or greater than the conventional limits of missing energy experiments such as ref. [3].

To summarize, a number of indications that dimension-3 SUSY breaking operators may not exist in the low energy effective theory were cited. We found that although gauginos are massless at tree level, radiative corrections give gluino masses in the $100-300$ MeV range and photino masses somewhat larger. The lightest R-hadron (the "glueballino", R^0) mass is estimated to be in the 1.4-2.2 GeV range, and its lifetime is likely to be longer than $\sim 10^{-10}$ sec. Therefore beam-dump experiments are more appropriately used to provide limits on squark masses than to exclude light gluinos. The scenario requires the mass of the lighter chargino to be below m_W, so it will be tested at LEP. Using signatures and detection strategies for R-hadrons and squarks developed in (II), positive evidence of this scenario could be found before that.

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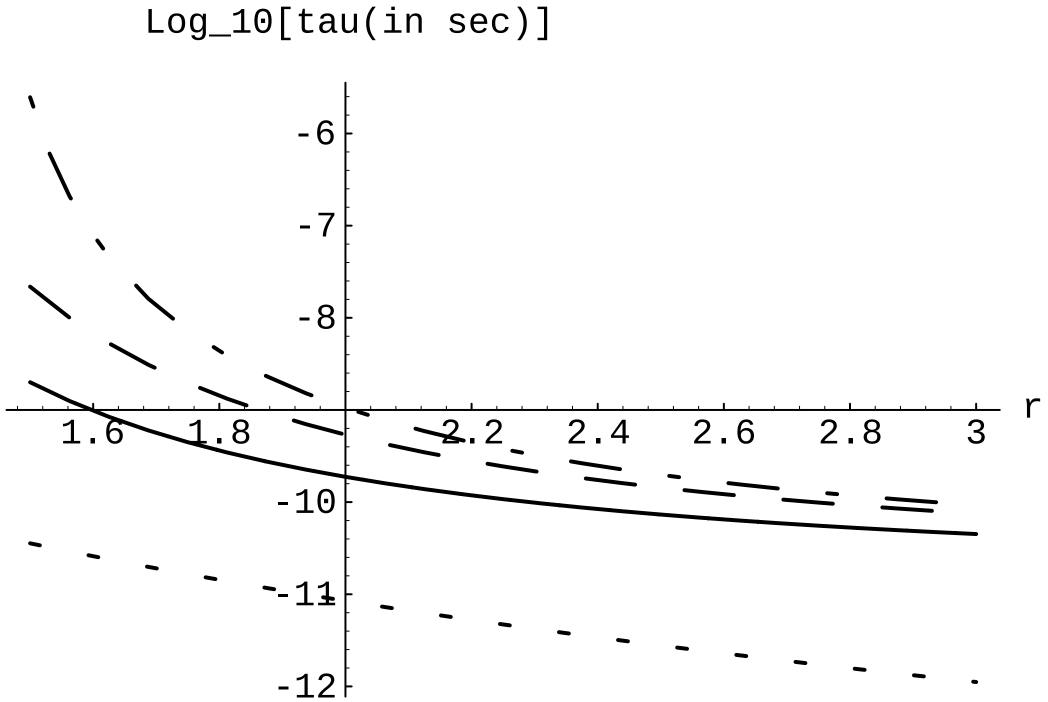

Figure 1: R^0 lifetime in a crude model for three different gluon distribution functions described in the text (solid: F_{ur}, dashed: F_{10}, dot-dashed: F_{nr}) as a function of $r \equiv \frac{m(R^0)}{m_{\tilde{\gamma}}}$, with $m(R^0) = 1.5$ GeV and $M_{sq} = 150$ GeV. The dotted curve is a plot of the lifetime of a free gluino of mass (r/1.5) GeV, decaying into massless $u\bar{u}$ or $d\bar{d}$ and$\tilde{\gamma}$ for $M_{sq} = 150$ GeV.

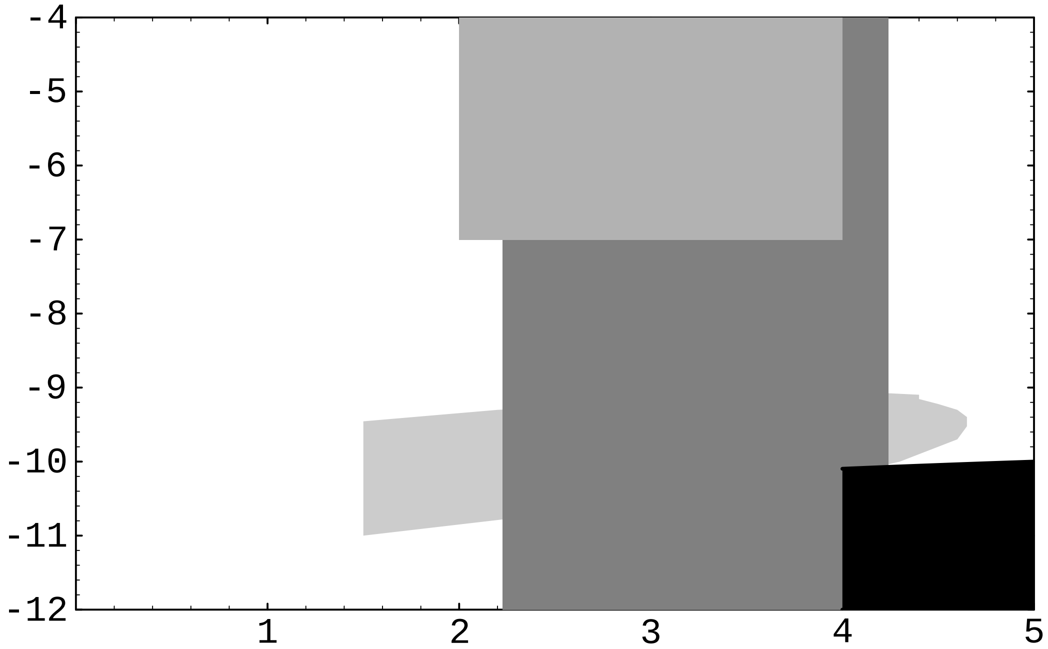

Figure 2: Experimentally excluded regions of $m(R^0)$ and $\tau_{\tilde{g}}$. Horizontal axis is $m(R^0)$ in GeV; vertical axis is Log_{10} of the lifetime in sec. A massless gluino would lead to $m(R^0) \sim 1.2 - 2.2$ GeV.