Superoblique corrections and nondecoupling of supersymmetry breaking

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If supersymmetric partners of the known particles have masses at the multi-TeV scale, they will not be directly discovered at planned future colliders and decouple from most observables. However, such superpartners also induce nondecoupling effects that break the supersymmetric equivalence of gauge boson couplings g_i and gaugino couplings h_i through supersymmetric analogues of the oblique corrections. Working within well-motivated theoretical frameworks, we find that multi-TeV scale supersymmetric particles produce deviations at the $1-10\%$ level in the ratios h_i/g_i. Such effects allow one to bound the scale of kinematically inaccessible superpartners through precision measurements of processes involving the accessible superparticles. Alternatively, if all superpartners are found, significant deviations imply the existence of highly split exotic supermultiplets. [S0556-2821(97)05223-5]

I. INTRODUCTION

Supersymmetric particles are often assumed to have a mass on the order of or below the TeV scale if supersymmetry (SUSY) indeed plays a role in the solution of the gauge hierarchy problem. Otherwise, fine-tuning of various parameters in the low-energy theory is required [1], undermining the motivation for the introduction of low-energy supersymmetry. The prospects for discovering and studying some supersymmetric particles (sparticles) at present and future colliders are therefore promising, particularly at the Large Hadron Collider (LHC) at CERN [2] and proposed high-energy linear e^+e^- colliders [3–5].

It is, however, a logical possibility that only part of the sparticle spectrum will be seen at planned future colliders, with some number of superpartners of ordinary matter and gauge fields beyond the discovery range. In fact, such scenarios are realized in a wide variety of models, and are often found in theories designed to solve the supersymmetric flavor problem, i.e., the problem that low-energy constraints are violated for generic sfermion masses and mixings. These models may be roughly divided into two categories. In the first class of models, which we will refer to as ''heavy QCD models,'' the gluino and all the squarks are heavy. Such may be the case in models with gauge-mediated SUSY breaking [6], where strongly interacting sparticles get large contributions to their masses, and also in the no-scale limit of supergravity models [7]. These models solve the SUSY flavor problem, since flavor-blind sfermion masses result from the proportionality of sfermion masses to gauge couplings and charges. A similar spectrum may also be predicted in other models, for example, grand-unified models with nonunified gaugino masses and heavy gluinos [8], in which the gluino drives the squark masses up through renormalization group evolution. In a second class of models, the first and second generation squarks and sleptons are very heavy with masses ~ 10 TeV, while the third generation sfermions are at the weak scale [9–12]. We will call these ''2–1 models.'' Such models are motivated by the desire to satisfy low-energy constraints from, for example, K^0-$\bar{K}^0$ mixing and $\mu \to e \gamma$, without the need for sfermion universality, sfermion alignment, or small CP-violating phases. At the same time, the extreme fine-tuning problem arising from very massive third generation sfermions is alleviated. It should be noted, however, that some increased level of fine-tuning must typically be tolerated, both in these models [9,13] and in those of the first category [14].

Given the many possibilities for supersymmetric particles beyond the reach of the LHC and proposed e^+e^- colliders, it is well worth considering what experimental implications such heavy states may have. In most experimentally accessible processes, such states decouple, and their effects rapidly decrease with increasing mass scale. Here, however, we study effects with the opposite behavior, that is, which grow with increasing supersymmetric mass splittings. Such effects rely on the fact that the interactions in supersymmetric theories are tightly constrained. For example, SUSY implies the relations

$$g_i = h_i, \qquad (1)$$

where g_i are the standard model gauge couplings, h_i are their supersymmetric analogues, the gaugino-fermion-sfermion couplings, and the subscript $i=1,2,3$ refers to the U(1), SU(2), and SU(3) gauge groups, respectively. Unlike other relations, such as the unification of gaugino masses, these relations hold in all supersymmetric models and are true to all orders in the limit of unbroken SUSY. However, SUSY-

breaking mass differences within superfields with standard model quantum numbers lead to corrections to Eq. (1) that grow logarithmically with the superpartner masses. Deviations from Eq. (1) are thus unambiguous signals of SUSY-breaking mass splittings, and by precisely measuring such deviations in processes involving accessible superparticles, bounds on the mass scale of the kinematically inaccessible sparticles may be determined.

The corrections to Eq. (1) are highly analogous to the oblique corrections [15,16] of the standard model. We will therefore refer to them as ''superoblique corrections'' and parametrize them by ''superoblique parameters,'' one for each gauge group. As is the case for oblique corrections, we will find that superoblique corrections are flavor independent and are enhanced for large heavy particle sectors. Furthermore, the simple nature of the corrections allows one to study them in a model-independent fashion using only TeV-scale parameters. As examples, we will calculate the size of these corrections in the two classes of models described above. In both cases, we find substantial contributions to all three superoblique parameters. Such corrections may be measured through a variety of processes, depending on what sparticles are available for study. Tests of the SU(2) relation $g_2 = h_2$ with charginos have been studied [17], as has the possibility of testing and looking for deviations in the U(1) relation with selectrons [18]. Soft SUSY-breaking effects on hard supersymmetric relations, i.e., relations between dimensionless couplings such as Eq. (1), were also noted in Ref. [19], where such effects were calculated for the specific case of squark widths. A general classification of possible observables at $e^{\pm} e^-$ and hadron colliders, as well as detailed studies of representative examples incorporating the variety of experimental uncertainties will be presented in an accompanying study [20].

We begin in Sec. II with a formal discussion of the flavor-universal corrections to Eq. (1). The analogy to the oblique corrections of the standard model is highlighted, and superoblique corrections and parameters are defined. In Sec. III flavor-dependent corrections, as well as other nondecoupling effects are discussed. In Sec. IV we estimate the size of the superoblique corrections in the heavy QCD and 2–1 models described above. The (typically small) contributions of vectorlike messenger and U(1)' sectors to these deviations are calculated in Sec. V. Our conclusions, as well as additional comments concerning possible implications of measuring superoblique corrections, are collected in Sec. VI.

II. SUPEROBLIQUE CORRECTIONS

We would like to identify robust experimental signatures of as-yet-undiscovered supersymmetric particles at future colliders. If only the standard model particles are available to us and we are only able to probe momentum scales below sparticle thresholds, broadly speaking, two approaches are possible. The first is to look for their virtual effects in low-energy processes. Unfortunately, in the models discussed in Sec. I with sparticle masses $\gtrsim 1 - 10$ TeV, such effects are often well below experimental sensitivities. This is just a statement of the decoupling theorem [21] for heavy super-partners from low-energy phenomena.

The second approach is to adopt some model-dependent

assumption such that the values of the low-energy couplings may be interpreted as signatures of heavy sparticles. For example, if one assumes grand-unification boundary conditions for the gauge coupling constants, their well-measured values at low energies are sensitive to sparticle thresholds. Threshold corrections have been extensively studied both with renormalization group techniques that incorporate leading logarithm effects [22] and through explicit one-loop calculations with finite corrections [23]. Such experimental signatures are, of course, model dependent and are obscured by other effects, such as grand-unified theory (GUT) scale threshold effects.[1]

In this study, we will consider scenarios in which some, but not all, superpartners are discovered. As noted in Sec. I, such scenarios may be realized at future colliders in a variety of models. If this is the case, what may be learned about the heavy, inaccessible superpartners? It is well known that the decoupling theorem does not apply if the heavy particle masses break symmetries [21]. In the present case, the heavy sparticle masses are predominantly invariant under standard model symmetries.[2] However, these masses violate supersymmetry, and in fact, the heavy superpartners give rise to nondecoupling corrections in processes involving the light superpartners. There are a variety of nondecoupling effects that may be considered. We will concentrate here on a set which we will call ''superoblique corrections,'' for reasons detailed below. These corrections are selected as particularly important, because they are universal in processes involving gauginos, enhanced by a number of factors, and may be measured at colliders in a variety of ways. Other nondecoupling effects will be described in Sec. III.

For simplicity, let us begin by neglecting the superpotential Yukawa couplings and assuming both R-parity (R_P) conservation [27] and flavor conservation. (The implications of relaxing these assumptions are the topic of the following section.) With these assumptions, in processes involving standard model particles or the light superpartners, the heavy superpartners appear at the one-loop level only through renormalizations of gauge boson and gaugino propagators. These renormalizations are equivalent in the limit of exact SUSY. However, since the sparticles have SUSY-breaking masses, the corrections from the heavy sparticle loops are different for gauge bosons and gauginos, and the effects are proportional to $\ln(M/m)$, where M (m) is the characteristic heavy (light) superpartner mass scale. These nondecoupling

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[1]Note, however, that it is possible that certain processes probe momentum scales above sparticle thresholds, even though no sparticles have been directly discovered. By extrapolating the low-energy couplings up to the characteristic momentum scales of such processes, the presence or absence of intermediate sparticle thresholds may be determined, independent of GUT assumptions. The possibility of such effects has been discussed, for example, in Refs. [24,25].

[2]Sfermion masses may break SU(2), but this breaking is typically suppressed by the left-right mixing $(m_{\text{fermion}}/m_{\text{sfermion}})^2$. Contributions of sfermions to the SU(2) oblique parameters therefore may usually be neglected [26], and are especially small in the scenarios we are considering, since the sfermion masses are at the multi-TeV scale.

effects are similar in origin to the logarithmically divergent loop corrections to the Higgs boson mass in supersymmetric theories [28]. In addition, they are process independent, up to small $O(p^2/M^2)$ corrections, where p is the momentum of the gauge bosons or gauginos, and can be absorbed into the gauge couplings g_i and gaugino couplings h_i.

It is instructive to draw an analogy between these effects and the oblique corrections [15,16] of the electroweak sector of the standard model.[3] In the standard model, heavy particles with isospin breaking masses enter low-energy observables dominantly through the vacuum polarization functions of the electroweak gauge bosons. More specifically, SU(2) multiplets with custodial SU(2) [30] breaking masses, such as the (t,b) multiplet, renormalize the propagators of the (W,Z) vector multiplet differently, leading to explicit custodial SU(2) breaking in the vector multiplet at the quantum level, and introducing nondecoupling effects that grow with the mass splitting. The supersymmetric nondecoupling corrections may be described analogously with the following replacements in the previous sentence: (i) SU(2) multiplets→supermultiplets; (ii) custodial SU(2) breaking masses→soft supersymmetry-breaking masses; (iii) (t,b) multiplet→$(\tilde{f},f)$ supermultiplet; (iv) (W,Z) vector multiplet→(gauge boson,gaugino) vector supermultiplet; (v) custodial SU(2) →supersymmetry. Motivated by the strength of this analogy, we will refer to the SUSY-breaking effects of the heavy superparticles as ''superoblique corrections.'' As is the case for the oblique corrections of the standard model, the superoblique corrections provide a unique opportunity to probe the scale of the heavy sector at low energies.

Let us investigate this analogy further. The oblique corrections of the standard model may be described in terms of the three parameters S, T, and U [15]. The latter two are measures of custodial isospin breaking, with the differences of the mass and wave function renormalizations of the W and Z (more correctly, W^3) at $p^2=0$ from heavy particles given by T and U, respectively. Below, we will define superoblique parameters that are measures of the splitting of g_i and h_i. Such coupling constant splittings are results of differences in the wave function renormalizations of gauge bosons and gauginos. The superoblique parameters we define are therefore most similar to U, and will be denoted by $\tilde{U}_i$, where the subscript i denotes the corresponding gauge group.

One might also hope that measureable supersymmetric analogues to S and T exist, especially since these are typically more sensitive probes of new physics in the standard model. The S parameter is a consequence of the extra U(1) gauge group, and is not a measure of custodial SU(2) breaking. There is therefore no analogous effect in supersymmetry. The analogue to T is a difference in the mass renormalizations of gauge bosons and gauginos. In our case, there is no mass renormalization of the gauge bosons due to the heavy superpartners if their masses are standard model gauge

invariant.[4] On the other hand, gaugino masses may receive contributions from heavy sparticle loops if these loops contain R-symmetry-breaking effects. If there were no tree level gaugino masses, or these masses were somehow known, the loop-generated gaugino masses [31] would be a probe of the mass splitting between components of a supermultiplet, providing a probe analogous to T, parametrized by three new superoblique parameters $\tilde{T}_i$. However, in a general softly broken SUSY theory, arbitrary gaugino masses already exist at tree level, and there is no tree level mass relation between the gauge bosons and the gauginos. [In contrast, custodial SU(2) symmetry enforces the relation $m_W=m_Z \cos\theta_W$ at tree level in the standard model.] The gaugino mass renormalizations therefore may be absorbed into these tree level terms, yielding no useful low-energy observables corresponding to T, unless one makes some assumptions about the tree level gaugino masses.

The nondecoupling SUSY-breaking effects may also be profitably understood in the language of renormalization group equations (RGE's). Above the heavy superpartner scale M, SUSY is not broken, and we have $h_i=g_i$. Below M, where the heavy superpartners decouple, light fermion loops still renormalize the gauge boson wave function (and thus, g_i) but heavy sfermion loops and sfermion-fermion loops decouple from gauge boson and gaugino wave function renormalization, respectively. (Gauge loops still renormalize both wave functions in the non-Abelian case.) Since not all loops from the supermultiplet decouple simultaneously, supersymmetry is broken in the gauge sector, and therefore the gauge couplings g_i and gaugino couplings h_i start to evolve differently.

The one-loop evolution of the gauge couplings between the heavy and the light superpartner scales gives

$$\frac{1}{g_i^2(m)} \approx \frac{1}{g_i^2(M)} + \frac{b_{g_i}}{8\pi^2} \ln\frac{M}{m}, \qquad (2)$$

where b_{g_i} is the one-loop β-function coefficient of the effective theory between the heavy and light mass scales, with the heavy superpartners decoupled. For the gaugino couplings, because SUSY is broken, the RGE's will depend on both gauge couplings g_i and gaugino couplings h_i. However, because the deviations of h_i from g_i are small, the contributions from this difference to the RG evolution are higher order effects and hence negligible. In addition, because $h_i \approx g_i$, the Ward and Slavnov-Taylor identities still hold approximately for the gaugino couplings, and the primary effect of the decoupled sparticles is to modify the one-loop β-function coefficient of the gaugino coupling RGE. Approximating $h_i \approx g_i$ in the RGE's, the gaugino couplings at the scales of the light and heavy sectors are thus related by

$$\frac{1}{h_i^2(m)} \approx \frac{1}{h_i^2(M)} + \frac{b_{h_i}}{8\pi^2} \ln\frac{M}{m}. \qquad (3)$$

The one-loop β-function coefficient b_{h_i} is obtained by subtracting the entire contribution of whole supermultiplets that

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[3]This analogy was previously noted by Randall [29].

[4]By mass renormalization here we mean the mass shift at $p^2=0$, i.e., the part which is independent of wave function renormalization. Note, however, that the physical masses m_W and m_Z are renormalized by wave function renormalization.

contain heavy superpartners. Substituting the supersymmetric boundary condition $g_i(M) = h_i(M)$, straightforward manipulations yield

$$\frac{h_i(m)}{g_i(m)} \approx 1 + \frac{g_i^2(m)}{16\pi^2}(b_{g_i} - b_{h_i})\ln\frac{M}{m}. \qquad (4)$$

To parametrize the nondecoupling effects of heavy superpartners, we define the superoblique parameters

$$\widetilde{U}_i \equiv \frac{h_i(m)}{g_i(m)} - 1 \approx \frac{g_i^2(m)}{16\pi^2}(b_{g_i} - b_{h_i})\ln R, \qquad (5)$$

where $i = 1,2,3$ denotes the gauge group, and $R = M/m$. As noted above, these parameters are supersymmetric analogues to the oblique parameter U [15], with one for each gauge group. Note that, because $b_{h_i} < b_{g_i}$, the coupling h_i is more asymptotically free than g_i, $h_i(m) > g_i(m)$, and the parameters $\widetilde{U}_i$ are always positive. This statement is true at the leading logarithm level irrespective of whether the heavy sparticles are scalars or fermions. We may also define another set of parameters that are deviations in the ratio of ratios, which we denote by the two-index variables

$$\widetilde{U}_{ij} \equiv \frac{h_i(m)/h_j(m)}{g_i(m)/g_j(m)} - 1 \approx \widetilde{U}_i - \widetilde{U}_j$$

$$\approx \frac{1}{16\pi^2}[g_i^2(m)(b_{g_i} - b_{h_i}) - g_j^2(m)(b_{g_j} - b_{h_j})]\ln R. \qquad (6)$$

These linear combinations of the superoblique couplings are useful, as they are probed by branching ratio measurements, which are sensitive to h_i/h_j.

In fact, the decoupling scales for the gauge and gaugino couplings are not identical when threshold corrections at the decoupling scale are taken into account. The finite threshold corrections slightly lower the decoupling scales for the gaugino couplings relative to those of the gauge couplings, which slightly reduces the deviations of h_i from g_i at low energy relative to the leading logarithm analysis. However, these effects may be absorbed into an effective heavy scale M', with $R = M'/m$. The finite corrections and the resulting shift in R are calculated in the Appendix.

III. OTHER NONDECOUPLING CORRECTIONS

In the discussion above, we have examined a set of nondecoupling corrections to the gaugino couplings that are universal in that they apply to all gaugino couplings. We have, however, neglected the superpotential Yukawa couplings and have also assumed R_P and flavor conservation. Such effects lead to additional nondecoupling corrections, including flavor-specific gaugino coupling corrections. In addition, couplings that do not involve gauginos also receive corrections (even in the absence of Yukawa couplings and R_P and flavor violation). Let us now consider each of these effects in turn.

In the presence of Yukawa couplings, new flavor-dependent nondecoupling radiative contributions are possible. For example, in the minimal supersymmetric standard model, matter field wave functions receive corrections from loops involving Higgs bosons and Higgsinos, Higgs and Higgsino wave functions are corrected by loops involving fermions and sfermions, and new contributions also appear in the vertices. These contributions grow logarithmically with the heavy mass in the loop. Such effects spoil the approximate Ward and Slavnov-Taylor identities for the gaugino couplings—if a gaugino coupling is renormalized by a Yukawa operator involving heavy superpartners, the diagrams involving the heavy field decouple and the cancellation of divergences is spoiled in the effective theory. The one-loop RGE's of the gaugino couplings will then also depend on Yukawa couplings, and the universal gaugino coupling h_i is split into different couplings h_i^f for each gaugino-f-$\tilde{f}$ vertex. These Yukawa coupling contributions to $\widetilde{U}_i$ are of the opposite sign to the universal corrections discussed above. Of course, such effects are typically suppressed by small Yukawa couplings and are only relevant for processes involving the Higgs, bottom, and top quark supermultiplets. Note that the RGE's now become dependent on all of the different gaugino couplings. (See, for example, Appendix B of Ref. [32].) However, such corrections from differences in the couplings are higher order effects, and may be neglected here.

An interesting case in which Yukawa couplings could be important is in theories with R_P violation. In the minimal supersymmetric standard model, lepton and baryon number are not accidental symmetries of the low-energy theory, but are put in by hand when one imposes R_P conservation. R_P-violating terms include Yukawa couplings of leptons λLLE, lepton doublets to quarks $\lambda' LQD$, and the different quark singlets $\lambda'' UDD$, where generational indices have been suppressed. Current bounds on individual couplings allow rather large couplings λ, λ', and λ'' for certain generational indices. (See, for example, Ref. [33], where present constraints on λ'_{333}, the only coupling with three third generation indices, are analyzed.) In addition, these bounds are often significantly weakened for heavy superpartner masses, and so, in the scenarios we are considering, may be extremely poor. Consequently, important negative and flavor-dependent Yukawa contributions to $\widetilde{U}_i^f$ could arise in R_P-violating models. Of course, R_P violation also allows the lightest supersymmetric particle to decay, leading to nonstandard supersymmetric signals, which modifies the strategies for measuring such superoblique parameters.

In the absence of flavor conservation, flavor mixing matrix elements will appear at the gaugino-fermion-sfermion vertices. In this case, if an sfermion in one generation belongs to the heavy sector and an sfermion in another generation belongs to the light sector, as may be the case, for example, in 2−1 models, heavy sfermion loops may appear in the matter wave function and vertex renormalizations of the gaugino couplings of the light sector through flavor-violating interactions. Such effects also contribute to the violation of the Ward identity for gaugino couplings. However, in such models, flavor mixings between the heavy and the light sectors are naturally suppressed by m/M. Therefore, the effects of these flavor-violating loop corrections should be small. Note, however, that such mixings may be measured or bounded by experiment [34], and such effects have implica-

tions for gaugino coupling measurements [20].

Up to this point, we have only discussed deviations of the SUSY relation between the gauge couplings and the gaugino couplings. In supersymmetric theories, there are also D-term quartic scalar couplings, which arise from SUSY gauge interactions, and are therefore proportional to g_i^2 in the SUSY limit. After the heavy superpartners decouple, the relations between the quartic scalar couplings and the gauge couplings also receive nondecoupling corrections (which can be viewed as superoblique corrections from the wave function renormalization of the auxiliary D fields), and also possibly the flavor-dependent corrections discussed above. However, such deviations are likely to be more difficult to investigate experimentally: the couplings of four physical scalars are extremely challenging to measure, and other probes of D terms, such as in Higgs boson decays and SU(2) doublet sfermion splitting, require ambitious measurements of other parameters, such as $\tan\beta$, the ratio of Higgs expectation values. Although such measurements may be possible in certain scenarios, in the rest of this study, we will concentrate on the superoblique corrections between the gauge couplings and gaugino couplings, which enter generically in all processes involving gauginos, and which appear much more promising experimentally.

IV. NUMERICAL ESTIMATES IN THEORETICAL FRAMEWORKS

In Sec. II, we discussed superoblique corrections in the general context of models with heavy and light sectors with arbitrary particle content. In this section, we will investigate what size corrections may be reasonably expected. For concreteness, we will consider the two well-motivated classes of models described in Sec. I, namely, ''heavy QCD models,'' in which the heavy sector includes all colored superpartners, and ''2–1 models,'' in which the heavy sector consists of the first two families of sfermions. We will estimate the contributions of the heavy sectors to the parameters $\widetilde{U}_i$ and $\widetilde{U}_{ij}$ in these two frameworks, treating all heavy sector particles as degenerate—nondegeneracies within the heavy sector typically only lead to higher order effects. Discussion of additional contributions to $\widetilde{U}_i$ in models that contain vector-like multiplets at some high mass scale, *e.g.*, in gauge mediation and U(1)′ models, is deferred to Sec. V. Note that while the results of this section are presented to serve as benchmarks, it is important to keep in mind that much larger effects may be possible from, for example, exotic particles.

A. Heavy QCD models

We first consider models with all strongly interacting sparticles in the heavy sector. This category includes models in which the sfermion and gaugino masses are dominated by a flavor-independent term that is a function of the low-energy gauge couplings. The hierarchy between the strong and electroweak gauge couplings is then translated into a mass hierarchy between colored and noncolored particles. Examples include the no-scale limit of minimal supergravity [7], in which scalar masses are determined only by gaugino loops, models with nonuniversal gaugino masses and a heavy gluino [8], in which squark masses are enhanced by gluino

TABLE I. The β-function coefficients and parameters $\widetilde{U}_i$ in the heavy QCD models. $R \equiv M/m$ is the ratio of heavy to light mass scales.

Gauge group G_i	b_{g_i}	b_{h_i}	$b_{g_i} - b_{h_i}$	$\widetilde{U}_i$
SU(3)	-7	Frozen	-7	$-6.7\% \times \ln R$
SU(2)	$-\frac{1}{2}$	$-\frac{7}{2}$	3	$0.80\% \times \ln R$
$\frac{5}{3}$U(1)	$\frac{11}{2}$	$\frac{33}{10}$	$\frac{11}{5}$	$0.29\% \times \ln R$

loop contributions, and gauge mediation models [6], in which the gaugino and sfermion masses are determined by gauge loops involving vectorlike messenger supermultiplets at the ~ 100 TeV scale.

In these models, minimization of the Higgs potential implies, given the constraint of the Z boson mass, that the Higgsino mass parameter μ is naturally of the order of the gluino mass. Thus, typically the Higgsinos and one Higgs doublet should be included in the heavy sector. However, the contributions of these particles to $\widetilde{U}_i$ and $\widetilde{U}_{ij}$ are small, and the primary impact of the scale of μ is on what experimental observables may be available to probe the superoblique corrections.

Assuming that the heavy sector consists of all squarks and the gluino, we present in Table I the β-function coefficients and the resulting parameters $\widetilde{U}_i$ from Eq. (5).[5] Inclusion of the Higgsinos and one Higgs doublet in the heavy sector would slightly enhance $\widetilde{U}_1$ and $\widetilde{U}_2$. We have chosen the grand-unification normalization for the hypercharge U(1); of course, $\widetilde{U}_1$ is independent of this choice. For simplicity, we assume $g_i(m) = g_i(M_Z)$ in our numerical estimates, which is sufficient for $m/M_Z \lesssim 3$. We also have

$$\widetilde{U}_{21} \approx \frac{1}{16\pi^2} \left[3g_2^2(m) - \frac{11}{5} g_1^2(m) \right] \ln R \approx 0.50\% \times \ln R. \tag{7}$$

The parameters $\widetilde{U}_i$ and $\widetilde{U}_{21}$ are logarithmically dependent on $R = M/m$; a typical value for this ratio in heavy QCD models is $R \sim 10$.

In these models, the gluino and the squarks are in the heavy sector and are decoupled. The coupling h_3 is therefore not renormalized below M, and by convention, we take its value below M to be frozen, with $h_3(Q<M) = h_3(M) = g_3(M)$. By assumption, the gluino and squarks are inaccessible at colliders, and so the parameter $\widetilde{U}_3$ may be measured only through their virtual effects. Such measurements are likely to be extremely difficult, as they require an understanding of process-dependent QCD corrections. Note, however, that if the gluino is light, then $b_{g_3} - b_{h_3} = 4$. Hence, the

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[5]The contribution of a (component) field j with spin S^j to the β-function coefficient b_i is $b_i^j = N_i^j a^j T_i^j$, where N_i^j is the appropriate multiplicity; $a^j = \frac{1}{3}, \frac{2}{3}, -\frac{11}{3}$ for $S^j = 0, \frac{1}{2}, 1$, respectively; and $T_i^j = 0, \frac{1}{2}, 2, 3$, or $\frac{3}{5} Y^2$ for a singlet, a particle in the fundamental representation of SU(N), an SU(2) triplet, an SU(3) octet, or, for $i = 1$, a particle with hypercharge $Y = I - Q$, respectively.

TABLE II. The parameter $\widetilde{U}_i$ in the 2–1 models.

Gauge group G_i	b_{g_i}	b_{h_i}	$b_{g_i} - b_{h_i}$	$\widetilde{U}_i$
SU(3)	$-\frac{13}{3}$	-7	$\frac{8}{3}$	$2.5\% \times \ln R$
SU(2)	$-\frac{1}{3}$	-3	$\frac{8}{3}$	$0.71\% \times \ln R$
$\frac{5}{3}$U(1)	$\frac{158}{30}$	$\frac{78}{30}$	$\frac{8}{3}$	$0.35\% \times \ln R$

sign of $\widetilde{U}_3$ could offer an indirect test of the ~ 1 GeV gluino scenario [35] if both gluinos and squarks are not observed at the LHC.

In the expressions above, we have treated all gaugino couplings as equivalent. In fact, as discussed in Sec. III, the various gaugino couplings may be significantly differentiated by Yukawa couplings. In this case, the gaugino-Higgsino-Higgs couplings h_i^H are split from the other gaugino couplings by nondecoupling corrections from the heavy t and b squarks. The corresponding parameters $\widetilde{U}_i^H$ are therefore diminished by the effects of the t- and b-quark Yukawa couplings and may be large and negative. For the remaining couplings, we have explicitly confirmed by comparison with the complete set of one-loop RGE's for heavy QCD models contained in Ref. [32] that the additional decoupling effects not included in Eq. (5) are negligible.

B. 2–1 models

Models with heavy first two generation scalars and light third generation scalars, Higgs fields, and gauginos have been discussed in Ref. [9], with explicit examples given recently in Refs. [10–12]. These models exploit and are motivated by the fact that the most stringent flavor-violating constraints may be satisfied by taking the sfermions of the first two families very heavy, while fine-tuning concerns may be alleviated by taking the other sparticles light, since the Higgs sector couples (at leading order) only to the sfermions of the third family and the electroweak gauginos. Note, however, that the heavy scale propagates to the light fields via hypercharge D terms and two-loop effects, leading to a strongly model-dependent upper limit on the heavy scale M [9,13]. Typical values of $R \sim 40 - 200$ may be taken in these models.

The gluino could, in principle, belong to either sector. For definiteness and motivated by gaugino mass unification, we will assume that all gauginos are in the light sector. The resulting parameters $\widetilde{U}_i$ are given in Table II. Since the decoupled sector consists of complete multiplets of a grand unified group, the differences $b_{g_i} - b_{h_i}$ are equal for all i, and the expressions for $\widetilde{U}_{ij}$ are simplified:

$$\widetilde{U}_{32} \approx \frac{1}{16\pi^2} \frac{8}{3} [g_3^2(m) - g_2^2(m)] \ln R \approx 1.8\% \times \ln R, \quad (8)$$

$$\widetilde{U}_{31} \approx \frac{1}{16\pi^2} \frac{8}{3} [g_3^2(m) - g_1^2(m)] \ln R \approx 2.2\% \times \ln R, \quad (9)$$

$$\widetilde{U}_{21} \approx \frac{1}{16\pi^2} \frac{8}{3} [g_2^2(m) - g_1^2(m)] \ln R \approx 0.35\% \times \ln R. \quad (10)$$

We see that the parameters $\widetilde{U}_3$, $\widetilde{U}_{32}$, and $\widetilde{U}_{31}$ are enhanced by the strong coupling and are therefore promising observables to probe.

Variants of 2–1 models may give alternative mass patterns, such as, for example, light and degenerate left-handed sleptons of the first two generations, and heavy right-handed selectrons and smuons [36]. A generalization of our results to these cases is straightforward. A reduced heavy sector diminishes $b_{g_i} - b_{h_i}$ and, thus, the corrections $\widetilde{U}_i$ and $\widetilde{U}_{ij}$. On the other hand, the existence of light selectrons and electron sneutrinos more than makes up for this setback, as it opens up the possibility of high precision probes of the electroweak superoblique parameters at $e^\pm e^-$ colliders that are inaccessible if these sleptons are all heavy [18,20].

V. VECTORLIKE (MESSENGER) SECTORS

The superoblique parameters receive contributions from all split supermultiplets with standard model quantum numbers. In many SUSY extensions of the standard model, there are extra vectorlike fields which transform under the standard model gauge groups. These vectorlike fields could have both SUSY-preserving and SUSY-breaking masses, and so they can also contribute to deviations in the SUSY relations $g_i = h_i$ at low energies. For example, this is the case in the gauge-mediated SUSY-breaking models, where the vectorlike messenger sector contains Dirac fermions with mass M_V and complex scalars with squared masses $M_V^2(1 \pm x)$. The low-energy ordinary sfermion spectrum is determined by M_V and x, and it is required that $|x| < 1$ in order to avoid tachyons and contradiction with experiments. More generally, irrespective of the mechanism that mediates SUSY breaking to the ordinary sector, there could exist some exotic vectorlike fields at or above the weak scale with Dirac fermions with mass M_V and complex scalars with squared masses $M_V^2(1 + x)$ and $M_V^2(1 + y)$. The variables x and y represent the SUSY breaking effects. If SUSY breaking is mediated through supergravity, x and y can be of order 1 only when the vectorlike fields have masses near the weak scale. If SUSY breaking is mediated through gauge interactions, x and y may be of order 1 only when the vectorlike fields are $\lesssim 100$ TeV; otherwise, through loop corrections, they will generate SUSY-breaking masses for standard model superpartners that are too large.

We consider first the messenger fields of gauge-mediated SUSY-breaking models. Let b_i be the contribution of the entire vectorlike supermultiplet sector to the appropriate one-loop β-function coefficient. For example, if the messenger sector contains n_5 pairs of **5** and $\overline{\mathbf{5}}$ and n_{10} pairs of **10** and $\overline{\mathbf{10}}$ SU(5) multiplets, then $b_i = n_5 + 3n_{10}$ for all i. If we naively perform a leading logarithm calculation, thereby ignoring finite pieces and decoupling all loops at the mass of the heaviest particle in the loop, we find

$$\delta\widetilde{U}_i^{\text{LL}} \approx \frac{g_i^2(M_V)}{64\pi^2} b_i \left(\frac{2}{3} \ln \sqrt{1+x} - \frac{1}{3} \ln \sqrt{1-x} \right). \quad (11)$$

As is evident from this expression, the leading logarithms $\ln(M_V/\mu)$ have canceled, as they must, since in this case, the SUSY breaking is governed not by M_V, but by x. The result is therefore reduced to a finite term, and we clearly must calculate the finite pieces correctly.

For gauge couplings, the naive decoupling is correct: the scalar loops decouple at $M_V\sqrt{1\pm x}$ and the fermion loops decouple at M_V. (See the Appendix.) The contribution to the gauge couplings can be written as

$$\delta\left(\frac{1}{g_i^2}\right) = -\frac{b_i}{8\pi^2}\frac{1}{2}\left[\frac{1}{3}\ln\frac{M_V\sqrt{1+x}}{\mu} + 2\times\frac{2}{3}\ln\frac{M_V}{\mu} + \frac{1}{3}\ln\frac{M_V\sqrt{1-x}}{\mu}\right]. \tag{12}$$

For the fermion-sfermion loop contribution to the gaugino wave function renormalization, we can apply the result in the Appendix. The Feynman parametrization integral of Eq. (A3) becomes

$$\int_0^1 d\alpha\, 2\alpha\, \ln\left[\frac{\alpha M_V^2(1\pm x)+(1-\alpha)M_V^2}{\mu^2}\right] = \ln\frac{M_V^2}{\mu^2} - \frac{1}{2}\pm\frac{1}{x} + \ln(1\pm x) - \frac{1}{x^2}\ln(1\pm x), \tag{13}$$

and so the contribution to the gaugino couplings is

$$\delta\left(\frac{1}{h_i^2}\right) = -\frac{b_i}{8\pi^2}\frac{1}{2}\left[\left(\ln\frac{M_V}{\mu} - \frac{1}{4} + \frac{1}{2x} + \frac{1}{2}\ln(1+x) - \frac{1}{2x^2}\ln(1+x)\right) + \left(\ln\frac{M_V}{\mu} - \frac{1}{4} - \frac{1}{2x} + \frac{1}{2}\ln(1-x) - \frac{1}{2x^2}\ln(1-x)\right)\right]. \tag{14}$$

As expected, we find that the $\ln(M_V/\mu)$ terms cancel in the difference between the g_i and h_i evolutions given in Eqs. (12) and (14), and the final result is

$$\delta\widetilde{U}_i = \frac{g_i^2(M_V)}{16\pi^2}b_i\left[-\frac{1}{4} + \left(\frac{1}{6} - \frac{1}{4x^2}\right)\ln(1-x^2)\right] \approx -\frac{g_i^2(M_V)}{384\pi^2}b_i x^2 \quad \text{for small } |x|. \tag{15}$$

The effect is very small for most of the range $0 \leq |x| < 1$, and it is therefore unlikely that any experimental measurement can be sensitive to superoblique corrections arising from such a messenger sector. Note, however, that this effect has a negative sign for small x relative to the logarithmic effect discussed in Secs. II and IV. Its smallness is thus fortunate, in the sense that such effects therefore cannot cancel the nondecoupling signatures of heavy superpartners.

It is also straightforward to obtain the result for the more general spectrum of vectorlike fields ($x \neq -y$):

$$\delta\widetilde{U}_i = \frac{g_i^2(M_V)}{16\pi^2}b_i\left[-\frac{1}{8} + \frac{1}{4x} + \left(\frac{1}{6} - \frac{1}{4x^2}\right)\ln(1+x)\right] + (x\to y) \approx \frac{g_i^2(M_V)}{16\pi^2}b_i\left(\frac{x+y}{12} - \frac{x^2+y^2}{48}\right) \quad \text{for small } |x| \text{ and } |y|. \tag{16}$$

For $x \neq -y$, the linear term does not vanish and we have a larger effect. However, unless there are many such heavy vectorlike multiplets (large b_i) with significant mass splittings among supermultiplet components (large $|x|$, $|y|$), the contributions to the superoblique corrections are small relative to the deviations discussed in Sec. IV. Note that in both the case of vectorlike messenger sectors and this more general case, large deviations are possible only for $|x|,|y| \approx 1$. If a deviation is seen which cannot be due simply to the standard model superpartners, the considerations stated above then strongly suggest that the masses of such vector-like particles are below the ~ 100 TeV scale.

VI. FINAL COMMENTS AND CONCLUSIONS

In this study we have considered low-energy softly broken supersymmetric theories that contain a heavy sparticle sector that is beyond the kinematical reach of planned future collider experiments. Sparticle spectra leading to such scenarios appear in certain limits of the most simple supergrav-

ity model, but more importantly, are known to arise in many other well-motivated frameworks for the soft SUSY-breaking parameters, and especially those that address the SUSY flavor problem. Here, we have shown that the heavy sparticle sector induces nondecoupling radiative corrections in the light sparticle sector, providing a crucial window for the exploration of the heavy sector through precision measurements in processes involving light sparticles.

The nondecoupling of SUSY breaking is analogous to the nondecoupling of SU(2) breaking in the standard model. Here we have considered a particularly important set of nondecoupling effects, which are analogous to the oblique corrections of the standard model, and which we therefore call superoblique corrections. Such corrections arise from gauge boson and gaugino wave function renormalization, and lead to deviations in the equivalence of gauge boson couplings g_i and gaugino couplings h_i. These corrections are therefore most closely identified with the oblique parameter U, and we have parametrized them with the superoblique parameters $\widetilde{U}_i \equiv h_i/g_i - 1$. The superoblique parameters have a number

of important features: they are model-independent measures of SUSY breaking, receive additive contributions from every split supermultiplet, and grow logarithmically with M/m, the ratio of heavy to light mass scales.

The superoblique parameters may be expressed simply in terms of $\ln(M/m)$ and group theory factors. As examples, we have estimated the corrections from heavy superpartners within specific theoretical frameworks and found typical values $\widetilde{U}_i \approx P_i \ln(M/m)$, where $P_i = 0.3\%$, 0.7%, 2.5% for $i = 1,2,3$, and the logarithm varies between 2 and 5. The hierarchy between the different parameters results from their proportionality to the low-energy gauge couplings, and the positive sign of the parameters is model independent at the leading logarithm level. We also calculated the contributions of messenger sectors in models of gauge mediation and possible exotic vectorlike multiplets. Such contributions were found to be typically very small, with substantial corrections only for highly split multiplets.

The effect of superoblique corrections in the accessible sparticles is to modify gaugino coupling constants. It is therefore not difficult to identify observables that are formally probes of such corrections. For example, the cross section of chargino production at e^+e^- colliders provides one such observable [17], as the gaugino couplings h_2 enter through t-channel sneutrino exchange. Selectron production at an e^+e^- collider provides another such probe [18]. In addition, if a particle has two or more decay modes, and at least one involves gauginos, its branching ratios are also probes of the superoblique corrections. Of course, all such measurements receive uncertainties from a variety of sources, ranging from backgrounds and finite statistics to the errors arising from the many other unknown SUSY parameters entering any given process. A classification of possible experimental probes at $e^{\pm}e^-$ and hadron colliders, as well as detailed studies of promising measurements incorporating such uncertainties, is contained in an accompanying article [20].

If superoblique corrections are measured, the implications are many and varied, depending on what precision is achieved and what scenario is realized in nature. The implications may be listed in increasing order of the precision of the measurements. If superoblique parameters are constrained to be roughly consistent with zero, such tests provide quantitative confirmation that such particles are indeed supersymmetric particles. If bounds on $\widetilde{U}_i$ at the level of $P_i\% \times \ln(M/m)$ are achieved, deviations from zero may be seen, providing evidence of a heavy sector. Finally, if bounds at the level of $P_i\%$ are achieved, the heavy mass scale may be constrained to within a factor of 3, providing a discriminant for model building, and in the most optimal scenarios, setting a target for future collider searches. Alternatively, if all superpartners are directly observed, deviations from $g_i = h_i$ are indications of the existence of, for example, exotic matter with highly split supermultiplets, which are likely to be below the ~ 100 TeV scale. If supersymmetry is discovered, the superoblique corrections will therefore provide a crucial window on the physics above the TeV scale.

Note added. While completing this work, we learned of related work in progress [37]. We thank D. Pierce, L. Randall, and S. Thomas for bringing this work to our attention.

APPENDIX: ONE-LOOP THRESHOLD CORRECTIONS AT THE HEAVY SUPERPARTNER MASS SCALE

In this appendix, we calculate the one-loop threshold corrections at the heavy superpartner scale. These finite corrections are usually included only when one uses two-loop RGE's. However, since in our case $\ln(M/m)$ is not necessarily very large, it is not clear *a priori* that the finite pieces are negligible relative to the leading logarithm contributions. It is therefore important that we consider these pieces in detail. This will be seen to be especially true when we consider the contributions from vectorlike messenger fields in models of gauge-mediated SUSY breaking, where the large logarithms cancel and the finite pieces must be treated carefully. This is discussed in Sec. V.

In calculating these corrections, we work in the SUSY preserving $\overline{\text{DR}}$ renormalization scheme, since we want to preserve the relation $g_i = h_i$ when SUSY is not broken.[6] The couplings measured at low energies should be converted into the same scheme before comparison.

We first consider the vacuum polarization of the gauge bosons due to the heavy scalar loops, $\Pi_H^{\mu\nu}(q) = (g^{\mu\nu}q^2 - q^\mu q^\nu)\Pi_H(q^2)$. The couplings are measured at much lower energies than the heavy scalar mass M_S, so we set the external momentum q to zero. The vacuum polarization is then given by the well-known result

$$
\begin{aligned}
\Pi_H(0) &= ig^2\mu^{4-d}T_R \int_0^1 d\alpha \int \frac{d^dk}{(2\pi)^d} \frac{(1-2\alpha)^2}{(k^2-M_S^2)^2} \\
&= -\frac{T_R}{3}\frac{g^2}{16\pi^2}\left(\frac{1}{2-d/2} - \gamma_E + \ln 4\pi - \ln\frac{M_S^2}{\mu^2}\right) \\
&\quad + O(4-d),
\end{aligned} \tag{A1}
$$

where here μ is the renormalization scale. T_R is defined by $T_R\delta^{ab} = \text{tr } T^aT^b$ and is $\frac{1}{2}$ for the fundamental representation of SU(N). We subtract the terms $1/(2-d/2) - \gamma_E + \ln 4\pi$ in the dimensional reduction with modified minimal subtraction ($\overline{\text{DR}}$) scheme. The remaining term, $\ln(M_S^2/\mu^2)$, vanishes when $\mu = M_S$, implying that the gauge coupling in the low-energy effective theory matches that in the high-energy theory at $\mu = M_S$. Therefore, we decouple the heavy scalar

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[6]In fact, our calculation is the same as in the modified minimal subtraction ($\overline{\text{MS}}$) scheme, as we only have scalars and fermions in the loop.

loops at the scale of their masses in calculating the low-energy gauge boson couplings. In doing so, there is no finite threshold correction at one loop.

Now we turn to the low-energy gaugino couplings. The

heavy loop of the gaugino wave function renormalization consists of a scalar and a fermion of masses M_S and m_f, respectively. The one-loop diagram gives

$$
\begin{aligned}
\Sigma_{2H}(q) &= i(-i\sqrt{2}h)^2\mu^{4-d}T_R\int\frac{d^dk}{(2\pi)^d}\frac{i(\slashed{k}+m_f)}{k^2-m_f^2}\frac{i}{(k-q)^2-M_S^2} \\
&= i2h^2\mu^{4-d}T_R\int_0^1 d\alpha\int\frac{d^dk}{(2\pi)^d}\frac{\alpha\slashed{q}+m_f}{[k^2+\alpha(1-\alpha)q^2-\alpha M_S^2-(1-\alpha)m_f^2]^2}.
\end{aligned}
\tag{A2}
$$

Setting the external momentum to zero, the contribution to the wave function renormalization is

$$
\begin{aligned}
\delta Z_2 &= \frac{d\Sigma_{2H}}{d\slashed{q}}\bigg|_{\slashed{q}\to 0} = i2h^2\mu^{4-d}T_R\int_0^1 d\alpha\int\frac{d^dk}{(2\pi)^d}\frac{\alpha}{[k^2-\alpha M_S^2-(1-\alpha)m_f^2]^2} \\
&= -T_R\frac{h^2}{16\pi^2}\int_0^1 d\alpha\, 2\alpha\left[\frac{1}{2-d/2}-\gamma_E+\ln 4\pi-\ln\frac{\alpha M_S^2+(1-\alpha)m_f^2}{\mu^2}\right]+O(4-d).
\end{aligned}
\tag{A3}
$$

For the fermion-sfermion loop, $m_f\simeq 0$, and the Feynman integral reduces to

$$
\int_0^1 d\alpha\, 2\alpha\,\ln\frac{\alpha M_S^2}{\mu^2}=\ln\frac{M_S^2}{\mu^2}-\frac{1}{2}=\ln\left(\frac{M_S e^{-1/4}}{\mu}\right)^2.
\tag{A4}
$$

In this case, there is a nonzero finite correction, which implies that the decoupling scale of the fermion-sfermion loop is at $M_S e^{-1/4}$ instead of M_S.[7] Therefore, to take account of the threshold corrections at the decoupling scale, we could replace the scale M in Eq. (3) by an effective decoupling scale $\widetilde{M}$ different from that in Eq. (2).

To get an understanding of how large such shifts in the decoupling scale are, let us consider theories with heavy sectors composed of scalars (and possibly gauginos) with mass M. Including the one-loop threshold corrections, we have

$$
\frac{1}{h_i^2(m)}\approx\frac{1}{h_i^2(M)}+\frac{b_i}{8\pi^2}\frac{1}{4}+\frac{b_{h_i}}{8\pi^2}\ln\frac{M}{m}-\frac{b_{h_i}}{8\pi^2}\frac{1}{4},
\tag{A5}
$$

where b_i is the one-loop β-function coefficient for both gauge and gaugino couplings above the squark scale. The deviation of h_i from g_i at low energies becomes

$$
\begin{aligned}
\frac{h_i(m)}{g_i(m)} &\approx 1+\frac{g_i^2(m)}{16\pi^2}(b_{g_i}-b_{h_i})\ln\frac{M}{m}-\frac{g_i^2(m)}{16\pi^2}(b_i-b_{h_i})\frac{1}{4} \\
&= 1+\frac{g_i^2(m)}{16\pi^2}(b_{g_i}-b_{h_i})\left(\ln\frac{M}{m}-\frac{3}{8}\right).
\end{aligned}
\tag{A6}
$$

Here we have used the relation $b_i-b_{h_i}=\frac{3}{2}(b_{g_i}-b_{h_i})$, valid since $b_i-b_{h_i}$ receives contributions from heavy scalars and their fermionic partners, while $b_{g_i}-b_{h_i}$ receives contributions only from the fermionic partners. We can see that the deviation is slightly smaller than that naively obtained by decoupling the heavy loop at the heaviest particle mass. However, as we are interested in the case where $\ln(M/m)\gtrsim 2$, the shift only introduces only a small correction to the total deviation.

---

[7]In the heavy Higgsino–light-Higgs-boson case, we have $\int_0^1 d\alpha\, 2\alpha\,\ln(1-\alpha)m_H^2/\mu^2=\ln(m_H^2/\mu^2)-\frac{3}{2}$.

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