Precision sparticle spectroscopy in the inclusive same-sign dilepton channel at LHC

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The inclusive same-sign dilepton channel is already recognized as a promising discovery signature for supersymmetry in the early days of the LHC. We point out that it can also be used for precision measurements of sparticle masses after the initial discovery. As an illustration, we consider the LM6 CMS study point in minimal supergravity, where the same-sign leptons most often result from chargino decays to sneutrinos. We discuss three different techniques for determining the chargino and sneutrino masses in an inclusive manner, i.e. using only the two well measured lepton momenta, while treating all other upstream objects in the event as a single entity of total transverse momentum $P_T$. This approach takes full advantage of the large production rates of colored superpartners, but does not rely on the poorly measured hadronic jets, and avoids any jet combinatorics problems. We discuss the anticipated precision of our methods in the early LHC data.

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A long standing problem in hadron collider phenomenology has been the determination of the absolute mass scale of new particles in events with missing energy. The prototypical example of this sort is provided by any model of low-energy supersymmetry (SUSY) with conserved $R$-parity, in which the lightest superpartner (LSP), typically the lightest neutralino $\tilde{\chi}_1^0$, is a neutral, weakly interacting particle of a priori unknown mass \cite{1}. Astrophysics also adds credence to such scenarios, since the LSP is a potential dark matter candidate, whose relic abundance is typically in the right ballpark \cite{2}. $R$-parity conservation guarantees that every event contains (at least) two invisible particles, whose energies and momenta are not measured, making the full reconstruction of such events a very challenging task.

Recently, several solutions to this problem at the Large Hadron Collider (LHC) have been proposed. Most of them rely on exclusive channels \cite{3}, where a sufficiently long decay chain can be properly identified. Unfortunately, this almost inevitably requires the use of hadronic jets in some form in the analysis – in most SUSY models, the main LHC signal is due to the strong production of colored superpartners, whose cascade decays to the neutral LSP necessarily involve hadronic jets. For many reasons, jets are notoriously difficult to deal with, especially in a hadron collider environment. Because of the high jet multiplicity in SUSY signal events, any jet-based analysis is bound to face a severe combinatorial problem and is unlikely to achieve any good precision. Thus it is imperative to have alternative methods which avoid the direct use of jets and instead rely only on the well measured momenta of any (isolated) leptons in the event.

In this letter, we describe three such methods, which are free of the jet combinatorial problem. For illustration, we shall use the standard example of $R$-parity conserving supersymmetry with a $\tilde{\chi}_1^0$ LSP. Its collider signatures have been extensively studied, and typically involve jets, leptons and missing transverse energy \cite{1}. Among those, the inclusive same-sign dilepton channel has already been identified as a unique opportunity for an early SUSY discovery at the LHC \cite{4,5}. The two leptons of the same charge can be easily triggered on, and provide a good handle for suppressing the SM background. In our analysis we use the LM6 CMS study point \cite{4}, whose relevant mass spectrum is given in Table 1. At point LM6, signal events with two isolated same-sign leptons typically arise from the SUSY event topology in Fig. 1. Consider the inclusive production of same-sign charginos, which decay leptonically as shown in the yellow-shaded box in the figure. The resulting sneutrino ($\tilde{\nu}$) could be the LSP itself, or, as in the case of LM6, may further decay invisibly to a neutrino $\nu$ and the true LSP $\tilde{\chi}_1^0$. Such same-sign chargino pairs typically result from squark decays, as indicated in Fig. 1. In turn, the squarks may be produced directly through a $t$-channel gluino exchange, or indirectly in gluino decays. Note that the two same-sign leptons in Fig. 1 are accompanied by a number of upstream objects (typically jets) which may originate from various sources, e.g. initial state radiation, squark decays, or decays of even heavier particles up the decay chain. In order to stay clear of jet combinatorial issues, we shall adopt a fully inclusive approach to the same-sign dilepton signature, by treating all the upstream objects within the black rectangular frame in Fig. 1 as a single entity of total transverse momentum $P_T$.

Given this very general setup, we now pose the following question: assuming that a SUSY discovery is made in the inclusive same-sign dilepton channel, is it possible to measure the individual sparticle masses $M_p$ and $M_c$ involved in the leptonic decays of Fig. 1, using only the

TABLE I: Selected sparticle masses (in GeV) at point LM6.

| $M_\tilde{q}_L$ | $M_{\tilde{\chi}^0_1}$ | $M_{\tilde{\chi}^+_1}$ | $M_{\tilde{\chi}^-_1}$ | $M_{\tilde{\nu}}$ | $M_{\tilde{\chi}_t}$ | $M_{\chi_0}$ |
|----------------|------------------------|------------------------|------------------------|----------------|------------------------|----------------|
| 939.8          | 339.3                  | 305.3                  | 291.0                  | 275.7         | 158.1                  |               |
transverse momenta of the two leptons $\vec{p}_{T}^{(1)}$ and $\vec{p}_{T}^{(2)}$, and the total upstream transverse momentum $\vec{P}_{T}$? Although it may appear that those three vectors do not provide a lot of information to go on, we shall show that this is possible. We discuss three different approaches.

**Method I.** Let us concentrate directly on the observed lepton momenta $\vec{p}_{T}^{(1)}$. Consider the two collinear momentum configurations illustrated in Fig. 2 and defined as follows. In each configuration, the lepton momenta are the same: $\vec{p}_{T}^{(1)} = \vec{p}_{T}^{(2)}$; and then they can be either parallel or anti-parallel to the measured upstream $\vec{P}_{T}$:

\[
s = +1 \Rightarrow \vec{p}_{T}^{(1)} = \vec{p}_{T}^{(2)} \uparrow \uparrow \vec{P}_{T};
\]

\[
s = -1 \Rightarrow \vec{p}_{T}^{(1)} = \vec{p}_{T}^{(2)} \downarrow \downarrow \vec{P}_{T}.
\]

In what follows we shall use the integer $s = +1 \text{ (} s = -1 \text{)}$ to refer to the parallel (anti-parallel) configuration: $s \equiv \cos(\vec{p}_{T}^{(1)}, \vec{P}_{T}) = \cos(\vec{p}_{T}^{(2)}, \vec{P}_{T})$. Now let us measure the maximum lepton momentum in each configuration:

\[
p_{CT}(sP_{T}) \equiv \max_{\vec{p}_{T}^{(1)} = \vec{p}_{T}^{(2)} \wedge \cos(\vec{p}_{T}^{(1)}, \vec{P}_{T}) = s} \left\{ \vec{p}_{T}^{(i)} \right\}.
\]

Observe that both $p_{CT}(+P_{T})$ and $p_{CT}(-P_{T})$ can be directly measured from the lepton $p_{CT}$ distributions. For example, construct a 2D scatter plot $(x, y)$ of

\[x = \cos(\vec{p}_{T}^{(1)} + \vec{p}_{T}^{(2)}, \vec{P}_{T}), \quad y = |\vec{p}_{T}^{(1)} + \vec{p}_{T}^{(2)}|,
\]

with the cut $|\vec{p}_{T}^{(1)} - \vec{p}_{T}^{(2)}| < \epsilon$ ($\sim 0$), and take the limit

\[
p_{CT}(sP_{T}) = \lim_{\epsilon \rightarrow 0} \left( \frac{y}{2} \right).
\]

Armed with the two measurements $p_{CT}(+P_{T})$ and $p_{CT}(-P_{T})$, we can now directly solve for the masses $M_{p}$ and $M_{c}$. The formula for $p_{CT}(sP_{T})$ is

\[
p_{CT}(sP_{T}) = \frac{M_{p}^{2} - M_{c}^{2}}{4M_{p}^{2}} \left( \sqrt{4M_{p}^{2} + (sP_{T})^{2}} - sP_{T} \right).
\]

Inverting (5), we get

\[
M_{p} = \sqrt{\frac{p_{CT}(-P_{T}) p_{CT}(+P_{T})}{p_{CT}(-P_{T}) - p_{CT}(+P_{T})}} P_{T},
\]
For us the importance of the $M_{T2}$ variable is that the momentum configurations in Fig. 2 are precisely the ones which determine its endpoint $M_{T2}^{max}$. The complete analytical dependence of the $M_{T2}$ endpoint $\hat{M}_p(\hat{M}_c, P_T)$ on both of its arguments $\hat{M}_c$ and $P_T$ is now known [7]:

$$\hat{M}_p(\hat{M}_c, P_T) = \begin{cases} \hat{M}_p(\hat{M}_c, +P_T), & \text{if } \hat{M}_c \leq M_c, \\ \hat{M}_p(\hat{M}_c, -P_T), & \text{if } \hat{M}_c \geq M_c, \end{cases}$$ (12)

where

$$\hat{M}_p(\hat{M}_c, sP_T) = \left\{ \begin{array}{l} p_{RT}(sP_T) \\ + \sqrt{\left(p_{RT}(sP_T) + \frac{sP_T}{2}\right)^2 + \hat{M}_c^2} \end{array} \right\} - \frac{(sP_T)^2}{4} \right\}^{\frac{1}{2}} .$$ (13)

Thus we can alternatively obtain the sparticle masses by measuring just two $M_{T2}$ kinematic endpoints, with arbitrary choices for the test mass $\hat{M}_c$ and the upstream $P_T$. For concreteness, let us pick some fixed $\hat{M}_c'$ and $P_T'$, and form the corresponding $M_{T2}$ distribution [9] and measure its endpoint $\hat{M}_p'$, also making a note of the configuration $s'$:

$$\left\{ \hat{M}_c', P_T' \right\} \xrightarrow{\text{measure}} \left\{ \hat{M}_p', s' \right\} .$$ (14)

Now perform a second such measurement

$$\left\{ \hat{M}_c'', P_T'' \right\} \xrightarrow{\text{measure}} \left\{ \hat{M}_p'', s'' \right\} .$$ (15)

By inverting (13), these two measurements allow the experimental determination of

$$p_{RT}(s'P_T') = \frac{\hat{M}_p'^2 - \hat{M}_c'^2}{4\hat{M}_p'^2} \left( 4\hat{M}_p'^2 + (s'P_T')^2 - s'P_T' \right)$$

and similarly for $p_{RT}(s''P_T'')$. Now taking the ratio

$$r = \frac{p_{RT}(s'P_T')}{p_{RT}(s''P_T'')} = \frac{\sqrt{4\hat{M}_p'^2 + (s'P_T')^2 - s'P_T'}}{\sqrt{4\hat{M}_p''^2 + (s''P_T'')^2 - s''P_T''}} ,$$ (17)

where in the second step we used eq. (9), we can solve (17) for the true parent mass $M_p$ in terms of measured quantities:

$$M_p = \begin{cases} -r s' P_T' s'' P_T'' \left( r - s' P_T' \right) \left( r - s'' P_T'' \right) \frac{1}{(1 - r^2)^2} , & \text{for } s' > s'', \\ -r s'' P_T'' s' P_T' \left( r - s'' P_T'' \right) \left( r - s' P_T' \right) \frac{1}{(1 - r^2)^2} , & \text{for } s' < s'', \end{cases}$$

and then find the true child mass $M_c$ from (9) as

$$M_c = M_p \left[ 1 - \left( 1 - \frac{\hat{M}_c^2}{M_p^2} \right) \frac{4\hat{M}_p'^2 + (s'P_T')^2 - s'P_T'}{4\hat{M}_p'^2 + (s''P_T'')^2 - s''P_T''} \right]^{\frac{1}{2}} ,$$ (19)

with $M_p$ already given by (18). Note than in this method, the values of $M_c', \hat{M}_p'', P_T'$ and $P_T''$ can be chosen at will, allowing for repeated measurements of $M_p$ and $M_c$.
and (19), as a function of the true input masses $M_\delta M$ and $\delta M_c$. FIG. 4: Scaling factors relating the error $M_\delta M$ and $\delta M_c$ to fitted masses $M_\ell$ and $M_p$, as a function of the true input masses $M_\ell$ and $M_p$.

derestimated, washing out the expected kink. There are two separate reasons behind this effect. Recall that the $M_{T2}$ endpoint on the left branch is obtained in the configuration $s = +1$ of Fig. 2 which requires the lepton to be emitted in the backward direction. As a result, the parent boost favors configurations with $s \simeq -1$ over $s \simeq +1$. Another consequence is that leptons with $s \simeq +1$ are softer and more easily rejected by the offline $p_T$ cuts. We conclude that $M_{T2}^{\text{max}}$ measurements on the left branch are in general not very reliable, and tend to jeopardize the traditional kink method. For example, using Method III to fit the data in Fig. 3 (green dotted line), we find best fit values of only $M_{p(fit)} = 212$ GeV and $M_{c(fit)} = 188$ GeV. Method I has a similar problem, since $p_T(\tau + p_T)$ is measured from events in the $s = +1$ configuration. Using the $M_p$ measurements from Fig. 3 at $M_c = 0$ and $M_p = 1$ TeV, we find from eq. (16) that $p_T(\tau + 420 \text{ GeV}) = 8.8$ GeV and $p_T(-420 \text{ GeV}) = 50.6$ GeV (compare to the nominal values of 14.8 GeV and 53.6 GeV, correspondingly). The resulting mass determination via eqs. (15) is $M_{p(fit)} = 212$ GeV and $M_{c(fit)} = 190$ GeV. We see that in both Method I and Method III, the masses are underestimated due to the systematic underestimation of the left $M_{T2}^{\text{max}}$ branch in Fig. 3. It is therefore of great interest to have an alternative method, which relies on the right $M_{T2}^{\text{max}}$ branch alone.

This is where the available freedom in Method II comes into play, since both test masses $M'_p$ and $M'_c$ can be chosen on the right branch. Taking $p_T^2 = 350 \pm 50$ GeV and $p_T^2 = 500 \pm 50$ GeV and repeating our earlier analysis, we find that $\delta M_p$ on the right branch is still on the order of 3 GeV, as in Fig. 3. The resulting error $\delta M_p$ on the measured parent (child) mass can be easily propagated from eqs. (18,19). Two ratios $\delta M_p/\delta M_c$ and $\delta M_c/\delta M_p$ are shown in Fig. 4 where for concreteness we have taken $M'_p = M'_c = 1000$ GeV. Fig. 4 reveals that the LM6 input values of $M_p$ and $M_c$ are rather unlucky, since the error $\delta M_p$ on the $M_{T2}$ endpoint is then amplified by a factor of almost 70. However, if $M_c$ and $M_p$ happened to be different, with the rest of the spectrum the same, the precision quickly improves. For example, with $\delta M_p = \pm 3$ GeV, the masses can be determined to within $\pm 30$ GeV ($\pm 75$ GeV) within the yellow (orange) region. One should keep in mind that the dominant uncertainty on $\delta M_p$ is due to the SUSY combinatorial background. We have verified that in the absence of such combinatorial background, $\delta M_p \lesssim 1$ GeV and the typical precision on $M_p$ and $M_c$ from Fig. 4 is then at the level of 10%.

In conclusion, we considered the inclusive same-sign dilepton channel in SUSY, which so far has only been used for discovery, but not for mass measurements. We demonstrated that it allows a separate determination of the chargino and sneutrino masses. We discussed three different methods, which rely exclusively on the well measured lepton momenta. The methods are completely general and inclusive, and can be applied to other SUSY topologies and to non-SUSY scenarios like UED.

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