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Constraining Light Colored Particles with Event Shapes

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Using recently developed techniques for computing event shapes with Soft-Collinear Effective Theory, LEP event shape data is used to derive strong model-independent bounds on new colored particles. In the effective field theory computation, colored particles contribute in loops not only to the running of α_s but also to the running of hard, jet and soft functions. Moreover, the differential distribution in the effective theory explicitly probes many energy scales, so event shapes have strong sensitivity to new particle thresholds. Using thrust data from ALEPH and OPAL, colored adjoint fermions (such as a gluino) below 51.0 GeV are ruled out to 95% confidence. This is nearly an order-of-magnitude improvement over the previous model-independent bound of 6.3 GeV.

Despite the fact that particle physics experiments have been running at and above 91 GeV center of mass energies for over two decades, it is not known if the standard model represents the complete particle content below this scale. For particles which carry no standard model quantum numbers, the only hope of producing them at colliders is through the Higgs, if there is a Higgs, and if they couple to it, or indirectly through off-shell intermediate states. But surprisingly, even colored particles, which interact with the strong force, are not significantly constrained. As long as they have small or vanishing couplings to electroweak gauge bosons, current data allows mass ranges well below the weak scale. A good example is a color adjoint Majorana fermion, such as the gluino in supersymmetric theories.

The gluino is a color octet and thus should have a large production cross section at hadron colliders, a non-negligible contribution to four-jet events at LEP, and a significant effect on the running of α_s . Most of the current bounds on a color octet fermion depend on how it hadronizes and how it decays. For example,

- If the gluino decays to two quarks and a very light neutralino, hadron collider data rules it out up to 308 GeV at 95% confidence level (C.L.) [1]. A recent study has shown that the Tevatron could probe gluino masses up to 150 GeV in the same decay channel independent of the neutralino mass [2].
- If the gluino is stable on detector lifetimes, ALEPH has excluded masses lighter than 26.9 GeV [3].
- A bound of 12 GeV, for a fixed $\alpha_s(m_Z) = 0.118$, has been set based on the gluino's potential contribution to the parton distribution functions [4]. A strict lower bound (*i.e.*, independent of α_s) has not been set.

As for a model independent limit, ALEPH [5] performed an analysis on four-jet observables as a measurement of the strong coupling constant and QCD color factors. The analysis found a good fit to QCD and ruled out gluinos below 6.3 GeV at 95% C.L.. An independent study [6], which included the use of electroweak precision data, arrived at the same lower limit. In both cases, the bound

comes essentially from the cross section for $q\bar{q}g\tilde{g}$ production which is very sensitive to the gluino mass at LEP 1 energies. The scale 6.3 GeV is where these searches lose sensitivity and can be taken as the current model-independent bound on the gluino mass.

Besides real production, new colored particles can be seen indirectly by their virtual effects. For example, any particle with color will contribute to the running of α_s . Some of the current model-independent bounds come from fitting the 1-loop β -function coefficient – which is sensitive to the number of flavors, n_f – to values of α_s measured at different energies. For example, DELPHI has done a study of mean values of event shapes and other inclusive observables leading to $n_f = 4.7 \pm 1.2$ (and $n_f = 4.75 \pm 0.44$ when combined with low energy thrust data) and ruling out gluinos less than 5 GeV. This study includes data from 14 to 200 GeV. However, by averaging the observables – for example, into the mean thrust $\overline{T}(Q)$ – this approach is not optimized to take full advantage of the available data.

A nominally positive feature of totally inclusive observables, such as $\overline{T}(Q)$, is that only one scale appears, so the perturbation series in α_s cannot be spoiled by the appearance of large logarithms. However, in searching for new particles through radiative corrections it is precisely these logarithms which have the most valuable information. In order to trust a differential calculation where the logarithms are relevant, the logs must be resummed. For many years resummation of event shapes was only available at next-to-leading order [7], which was insufficient to provide strong bounds on new physics because of large theoretical uncertainties. Recently, however, the thrust distribution was resummed to next-to-next-to-next-to-leading logarithmic order using techniques of effective field theory [8]. Including matching to recent next-to-next-to-leading fixed-order (NNLO) event shapes [9], the theoretical uncertainty on the α_s extraction from LEP was finally reduced to be sub-dominant to other uncertainties for the first time. Moreover, besides reducing the uncertainty, the effective theory approach makes explicit that α_s is probed at many scales and so the sensitivity to new physics should be strong. Thus, it is natural to

try to improve the model-independent bounds on new colored states using these recent theoretical advances. In this letter, we use these insights to improve the model-independent bound on the gluino by nearly an order of magnitude.

The thrust distribution was shown in [8, 10] to have the form

$$R(\tau) = \frac{1}{\sigma_{\text{had}}} \int_0^\tau \frac{d\sigma}{d\tau'} d\tau' = \frac{1}{\sigma_{\text{had}}} [R_2(\tau) + r(\tau)] \quad (1)$$

where $\tau \equiv 1 - T$. Here, the matching function $r(\tau)$ is defined as the difference between the fixed-order thrust distribution and the fixed-order expansion of the resummed distribution. We use $r(\tau)$ at next-to-next-to-leading order, *i.e.* to α_s^3 .

The function $R_2(\tau)$ in Eq.(1) is the resummed distribution. It can be calculated using Soft-Collinear Effective Theory [11, 12, 13] with insights from [14, 15, 16, 17, 18, 19]. The result is [8]:

$$\begin{aligned} R_2(\tau) = & \exp [4S(\mu_h, \mu) - 2A_H(\mu_h, \mu) - 8S(\mu_j, \mu) \\ & + 4A_J(\mu_j, \mu) + 4S(\mu_s, \mu) + 2A_S(\mu_s, \mu)] \\ & \times H(Q, \mu_h) \left[\tilde{j}(\partial_\eta, \mu_j) \right]^2 \tilde{s}_T(\partial_\eta, \mu_s) \left[\frac{e^{-\gamma_E \eta}}{\Gamma(\eta + 1)} \right]. \quad (2) \end{aligned}$$

The derivatives in Eq. (2) are to be taken analytically and then η set to its canonical value $\eta = 4A_\Gamma(\mu_j, \mu_s)$. Here, $H(Q, \mu)$ is the hard function and $\tilde{j}(L, \mu)$ and $\tilde{s}(L, \mu)$ are the Laplace transforms of the jet and soft functions; all of these have power series expansions. $S(\nu, \mu)$ and $A_X(\nu, \mu)$ are auxiliary functions defined as integrals over various anomalous dimensions. Explicit expressions for these functions can be found in [8]. The scale μ in Eq.(2) is arbitrary and the distribution is formally independent of it, but different values of μ can be chosen for calculational convenience.

The formula (2) is a simplified version of the one in [8], valid when the scales are set to their canonical values:

$$\mu_h = Q, \quad \mu_j = Q\sqrt{\tau}, \quad \mu_s = Q\tau \quad (3)$$

As mentioned above, a calculation of mean thrust, or a fixed-order calculation of differential thrust, would only probe $\alpha_s(Q)$ at a single scale, the hard scale $\mu_h = Q$. But the differential thrust distribution probes even lower scales. For example, in the two-jet limit thrust reduces to the sum of hemisphere masses, $Q^2\tau \sim M_L^2 + M_R^2$. The effective theory expression associates this mass scale with the scale of jet functions, and probes it through $\mu_j \sim Q\sqrt{\tau}$. Actually, the effective theory makes it apparent that even lower scales, associated with soft modes of QCD, are relevant. These are probed by the soft scale $\mu_s \sim \mu_j^2/\mu_h \sim Q\tau$, which is a type of seesaw scale lower than both of the physical external scales Q and $Q\sqrt{\tau}$ [10]. Since α_s is larger at lower energy, resumming logs of the

soft scale is critical to generating an accurate thrust distribution.

Because the differential thrust distribution is sensitive to many scales, it would be sensitive to the presence of new colored particles with a variety of masses. These new states would affect the running of α_s , through the QCD beta function, as well as the hard, jet, soft anomalous dimensions – which appear implicitly in (2) – and the fixed-order hard, jet, and soft functions, H, \tilde{j} and \tilde{s} .

Throughout the following we modify the standard model by adding Δn_f new flavors of mass m at a threshold scale μ_{th} . For example, a new massive quark corresponds to $\Delta n_f = 1$ and a gluino to $\Delta n_f = 3$ [20]. Below the scale μ_{th} , the new flavors are integrated out, insuring decoupling as $m \rightarrow \infty$. This will, in general, induce discontinuities in $\alpha_s(\mu)$ and in the hard, jet and soft functions, all of which are unphysical by themselves. The resulting thrust distribution, however, must be smooth. In fact, one can show that the effective field theory distribution is independent of μ_{th} order-by-order in perturbation theory [21]. For simplicity, we take $\mu_{\text{th}} = m$ and match α_s at one-loop. To avoid having to run the jet and soft functions through the threshold, we choose $\mu = m$ in Eq.(2) when $\mu_s < m < \mu_j$. For $m < \mu_s$, we take $\mu = \mu_s$ and for $m > \mu_j$ we take $\mu = \mu_j$.

To demonstrate the sensitivity to new states, we begin by looking at a single data set, the ALEPH data from LEP 1 at 91.2 GeV [22]. We perform a bin-by-bin correction for hadronization and quark masses using PYTHIA v. 6.409 [23]. Using the fit region $0.10 \leq \tau \leq 0.24$ the soft scale μ_s probes 9 - 22 GeV and the jet scale μ_j probes 29 - 45 GeV. Thus, if there are n_f flavors below 9 GeV, we do not have to worry about an explicit threshold and may simply run α_s and the other objects using this value of n_f throughout (a more refined procedure is described below). To derive a bound on the number of light flavors, we perform a least-squares fit to the experimental data. For the errors used in the fit, we include both the experimental statistical uncertainty and also the statistical uncertainty in the fixed-order thrust distribution. The fixed-order result was calculated numerically, with somewhat slow convergence at NNLO, and to be conservative we rescale the NNLO uncertainties by a factor of 1.5 to account for the fact that the errors may have been underestimated. A combined fit with two free parameters gives $\alpha_s(m_Z) = 0.1169 \pm 0.0004$ and $n_f = 5.32 \pm 0.59$, where the errors are statistical only.

For a second example, using the same LEP 1 data set, we note that a gluino of mass $m = 25$ GeV lies outside the range of scales probed by the hard, jet, and soft functions. Thus, it can be modeled by taking $\Delta n_f = 3$ for the hard and jet functions, and $\Delta n_f = 0$ for the soft function. Performing the fit with these values, we find $\chi^2 = 31.7$ with the gluino compared to $\chi^2 = 11.9$ for the standard model, with 13 degrees of freedom. The fits for the two

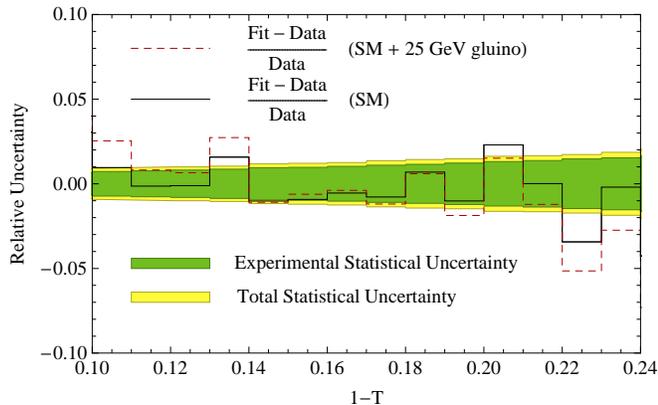


FIG. 1: Theoretical prediction versus ALEPH data at LEP 1 for the standard model and the standard model with a 25 GeV gluino. The total statistical uncertainty band includes theoretical statistical uncertainty from the Monte Carlo used to generate the NNLO fixed-order thrust distribution.

models are shown in Figure 1, where it is clear that the model with the gluino is systematically worse.

To properly scan over masses, we must specify how to handle the thresholds. First, consider the total hadronic cross section, σ_{had} . The exact leading order dependence of σ_{had} on the new particle mass can be extracted from [28]. For $m < \mu$, the contribution to the total cross section is proportional to $\Delta\sigma_{\text{had}} = \alpha_s^2(\mu) \left(\rho_V \left(\frac{m^2}{Q^2} \right) + \rho_R \left(\frac{m^2}{Q^2} \right) + \frac{1}{4} \log \left(\frac{m^2}{\mu^2} \right) \right)$, where ρ_V is the virtual contribution which vanishes at $m = \infty$ and ρ_R is the real emission contribution which vanishes for $m > Q/2$. The explicit log compensates the μ -dependence of α_s and is necessary to have a smooth $m \rightarrow 0$ limit. We will use this exact expression $\Delta\sigma_{\text{had}}$ for the new physics contribution to σ_{had} in Eq. (1), but observe that, as shown in [28], it is well approximated for $0 < m < Q$ by the leading power in m^2/Q^2 . Actually, it is not clear whether the experiments would have included decay products of real gluinos in their event selection for the thrust distribution, so in the spirit of providing a model-independent bound, we allow $\Delta\sigma_{\text{had}}$ to scan between 0 and the cross section for Δn_f additional massless flavors. This variation is included in the uncertainty band described below.

The exact contributions of massive colored states to the jet, soft, and hard functions are not known, but since the same loops and real-emission diagrams are relevant for them as for $\Delta\sigma_{\text{had}}$, it is likely that the result would be similar to that of $\Delta\sigma_{\text{had}}$. Thus, we assume the leading power is linear in m^2/μ_n^2 for the hard function, m^2/μ_j^2 for the jet function, and m^2/μ_s^2 for the soft function. That is, we take H, \tilde{j} and \tilde{s} to interpolate between the expression for $n_f = 5 + \Delta n_f$ flavors at $m = 0$ and $n_f = 5$ flavors at the relevant threshold. This removes any remaining discontinuity in the thrust distribution, and should be a

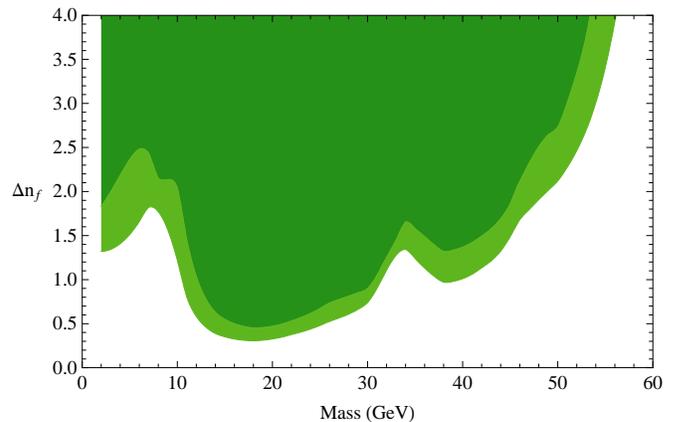


FIG. 2: Bounds on light colored particles from LEP data. The darker region is completely excluded at 95% confidence. The lighter region is an uncertainty band including estimates of various theoretical uncertainties.

good approximation to the (unknown) exact result. In a similar vein, the matching correction, $r(\tau)$ in Eq. (2), formally takes place at the hard scale Q . However, it depends on n_f and would be discontinuous as m crosses Q unless the discontinuity is removed by inclusion of explicit mass corrections. We use an interpolation also linear in m^2/Q^2 for this effect. Using this model for the mass thresholds, in lieu of the exact result, introduces some theoretical uncertainty. To account for that uncertainty, we explore some variations of the model and include the errors in our final bound, as described below.

With this treatment of the threshold effects, the thrust distribution is smooth and can be compared with the data for each m and Δn_f . We perform a combined fit to the ALEPH [22] and OPAL [24, 25] data sets from 91.2–206 GeV [26, 27]. The fit regions used are $0.1 < \tau < 0.24$ for LEP 1, and $0.04 < \tau < 0.25$ for ALEPH LEP 2 and $0.05 < \tau < 0.22$ for OPAL LEP 2. The data are corrected bin-by-bin for hadronization and bottom/charm mass effects using PYTHIA. We perform a least-squares fit of the theoretical prediction to the corrected data, using errors which include both the experimental statistical errors and the statistical errors of the NNLO fixed-order calculation, rescaled by 1.5, as described above. For the standard model, the χ^2 is 85.7 for 78 degrees-of-freedom. For each value of m and Δn_f , we minimize χ^2 and compute the maximum likelihood ratio as compared with the standard model. The resulting 95% C.L. bound is shown in Figure 2. For $\Delta n_f = 3$, the limit is $m_{\tilde{g}} > 52.5$ GeV. For a real gluino (with the appropriate group theory factors differing from $\Delta n_f = 3$ at higher orders), the bound differs by 0.03 GeV.

To account for the theoretical uncertainty, we include an uncertainty band (the light shaded region in Figure 2). This subsumes the following variations: (i) Removing the lowest bins from each data set in the fit. (ii) Not

interpolating the total cross-section and matching correction (we include variations both with $n_f = 5$ and with $n_f = 5 + \Delta n_f$ in σ_{had} and $r(\tau)$). (iii) Varying parameters in the hadronization model between PYTHIA's default values and the ALEPH [29] and OPAL [30] optimized tunings (PARJ(81) = 0.290, 0.292, 0.250 and PARJ(82)=1.00, 1.57, 1.90). (iv) Using a power correction proportional to m/μ_x instead of m^2/μ_x^2 for the threshold corrections. The band in Figure 2 includes the maximal and minimal bounds at 95% C.L. for each value of m and Δn_f . For the gluino, this gives $m_{\tilde{g}} > 51.0 - 54.0$ GeV. For our final bound, we take the least restrictive value, $m_{\tilde{g}} > 51.0$ GeV.

From Figure 2, it is clear that this method has the strongest sensitivity in an intermediate mass range, 10 GeV $\lesssim m \lesssim 40$ GeV. This range roughly coincides with the scales probed by the jet and soft functions in the fit regions of the thrust distributions. For masses below about 10 GeV, the mass threshold lies outside of the fit regions and the effect on the event shape can be partially compensated by a change in α_s . When the mass falls inside the range of the thrust distribution, it is more difficult to compensate by rescaling α_s , hence the stronger bound. With additional independent constraints on α_s , for example, from the lattice [31] or from τ decays [32], one might be able to close the light mass window more tightly. This might, for example, even rule out additional light colored triplets or scalar adjoints. However, as the lattice and τ -decay determinations of α_s (which take place at similar scales) are themselves inconsistent by more than two standard deviations, it is unclear whether a definitive bound could be obtained in this way. The main result of this paper is that event shapes alone are sufficient to exclude light and intermediate mass gluinos up to 51 GeV, independently of their decays.

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