Optical control methods offer tantalizing possibilities for creating “designer” two-body interactions in ultracold atomic gases with both high spatial resolution and high temporal resolution. By controlling the elastic scattering length, the inelastic scattering length, and the effective range, optical methods enable control of few-body and many-body systems, opening new fields of research.

Single optical field methods have been used to control Feshbach resonances, with large detunings to suppress spontaneous scattering, leading to limited tunability [1–9]. The first experiments, by Lett and collaborators in a Na Bose gas, created optical Feshbach resonances by coupling the ground and excited molecular states in the input channel [3]. Recently, Chin and co-workers observed suppression of both spontaneous scattering and the polarizability of Cs atoms by tuning between the $D1$ and $D2$ lines [9]. This method suppresses unwanted optical forces and achieves a lifetime of up to 100 ms with a rapid tuning of up to tens of milligauss by modulating the intensity of the control beam. Cetina and co-workers [10] vary the intensity of a far-detuned optical trapping potential to achieve a 40 mG tuning, arising from the unequal light shifts of atomic and molecular states.

Building on ideas suggested by Bauer et al. [7] and Thalhammer et al. [11], we are developing two-field optical control methods [12,13] that create a molecular dark state in the closed channel of a magnetic Feshbach resonance, as studied recently in dark-state spectroscopy [14]. This approach is closely related to electromagnetically induced transparency (EIT) [15], where quantum interference suppresses unwanted optical scattering. Our two-field methods are particularly important for broad Feshbach resonances with large background scattering lengths, where they enable symmetrical tuning of the scattering length about the minimum loss point, with a much smaller loss than single-field methods for the same tuning [16]. The two-field methods also produce narrow energy-dependent features in the scattering phase shift, enabling large changes in the scattering length with small changes in the optical frequencies compared to single-field methods [16]. By using the optical frequency rather than the intensity as the control parameter, unwanted changes in the total trapping potential are suppressed. Furthermore, these narrow features enable control of the effective range [13]. Analogous to the EIT method of enhancing optical dispersion in gases with suppressed absorption [15], the effective range can be modified in regions of highly suppressed optical scattering [16].

Implementation of optical control methods requires an understanding of the optically induced level structure and energy shifts, which depend on the relative momentum of a colliding atom pair. Our original theoretical approach [12,13] and that of other groups [7] employed adiabatic elimination of an excited molecular state amplitude, which fails for very broad resonances where the hyperfine coupling constant is large compared to the radiative decay rate. This unresolved issue has been noted previously [21].

In this Letter, we demonstrate large shifts of magnetic Feshbach resonances and strong suppression of spontaneous scattering in measurements of two-field light-induced loss spectra. Furthermore, we present a new theoretical approach to describe control of broad and narrow Feshbach resonances in a unified manner, replacing a “bare”-state description by a more natural description in terms of “continuum-dressed” states that incorporate the hyperfine coupling into the basis states. Using the measured Rabi frequencies, the predicted relative-momentum averaged loss spectra agree in shape and magnitude with data for both broad and narrow resonances.
Excited vibrational state, which decays at a rate \( \gamma_e = 2\pi \times 11.8 \text{ MHz} \).

The detunings of the optical fields that couple state \(|g_1\rangle\) and \(|g_2\rangle\) to \(|e\rangle\) are \( \Delta_1 \) and \( \Delta_2 \), respectively. The single photon detuning of the \( \omega_1 \) beam for the \( |T\rangle \rightarrow |e\rangle \) transition is a function of the magnetic field and can be defined at a reference magnetic field \( B_{\text{ref}} \) as \( \Delta_0 = \Delta_L - 2\mu_B (B - B_{\text{ref}})/\hbar \), where \( \Delta_L \) is the detuning of the optical field when \( B = B_{\text{ref}} \). The two-photon detuning for the \( |T\rangle \rightarrow |e\rangle \rightarrow |g_2\rangle \) system is \( \delta = \Delta_e - \Delta_2 \).

We prepare a 50–50 mixture of \(^{6}\text{Li}\) atoms in the two lowest hyperfine levels, \( |1\rangle \) and \( |2\rangle \), in a CO\(_2\) laser trap with trap frequencies \( \omega_x \approx 2\pi \times (3100, 3350, 120) \text{ Hz} \). The \( |1\rangle \rightarrow |2\rangle \) mixture of \(^{6}\text{Li}\) has a broad Feshbach resonance at \( B_{\text{ref}} = 832.2 \text{ G} \) [22,23] of width \( \Delta B_B = 500 \text{ G} \) due to strong hyperfine coupling of the triplet continuum to the “broad” singlet state \(|g_1\rangle_B\) [13]. In addition, there is a narrow Feshbach resonance at 543.2 G of width \( \Delta B_N = 0.1 \text{ G} \) [24] due to weak second order hyperfine coupling of the triplet continuum \(|T\rangle\) to the narrow singlet state \(|g_1\rangle_N\) [13]. After forced evaporation, and reraising the trap to full trap depth, we have approximately \( 10^8 \) atoms per spin state for our experiments. We generate both the \( \omega_1 \) and \( \omega_2 \) beams from diode lasers, locked to a stabilized cavity near 673.2 nm. The relative frequency is \( \approx 57 \text{ GHz} \) [14] with a jitter \( < 100 \text{ kHz} \). The absolute frequency stability is \( < 100 \text{ kHz} \). Both laser beams are polarized along the bias magnetic field \( z \) axis [16].

Initially, we use a single field \( (\omega_1) \) to observe the shift of the narrow Feshbach resonance in the atom loss spectra. After forced evaporation at 300 G, the magnetic field is swept to the field of interest and allowed to stabilize for \( \approx 1 \text{ s} \), which produces a background three-body loss at 543.2 G. Then the \( \omega_1 \) field illuminates the atoms for 5 ms, with a detuning \( \Delta_L = 30.2 \text{ MHz} \) (with \( \Delta_L \equiv 0 \) for \( B = B_{\text{ref}} = 543.2 \text{ G} \)), after which the atoms are imaged at the field of interest to determine the density profile and the atom number.

The single-field atom loss spectra versus the magnetic field (see Fig. 2) exhibit two loss peaks. (i) A broad peak
arises at 554 G, where the \( \omega_1 \) optical field is resonant (\( \Delta_e = 0 \)) with the \( |T\rangle \rightarrow |e\rangle \) transition. Here, the transition arises from the hyperfine coupling of \( |T\rangle \) to \( |g_1\rangle_B \), far from the resonance at 832.2 G. (ii) A narrow peak below 543.2 G occurs as the magnetic field tunes the triplet continuum near \( |g_1\rangle_N \), which is light shifted in energy—and hence, in the magnetic field, to \( B'_{\text{res}} \)—due to the \( \omega_1 \) optical field, detuned from the \( |g_1\rangle_N \rightarrow |e\rangle \) transition by \( \Delta_1 = \Delta_e = 30.2 \) MHz (see Fig. 1). At this field, the transition strength is maximized, while the \( \omega_1 \) optical field is off resonant with the \( |T\rangle \rightarrow |e\rangle \) transition by \( \Delta_1 = 30.2 \) MHz. For a \( \Omega_1 = 3.1\gamma_e \), the narrow resonance is shifted downward by 3.0 G; approximately 30 times the width \( \Delta B_N \). The continuum-dressed-state model (the solid red line) reproduces the shift of the narrow resonances and the amplitudes of both the narrow and broad resonances, using the measured \( \Omega_1 \) [16].

We suppress loss by applying a second field \( \omega_2 \) (see Fig. 3). At the magnetic field of interest, the \( \omega_2 \) beam is adiabatically turned on over 30 ms. Then the \( \omega_1 \) beam is applied for 5 ms, after which both beams are turned off abruptly. The \( \omega_2 \) beam creates an optical dipole trap and provides additional confinement in the \( z \) direction, due to its high intensity, changing the axial trap frequency from 120 Hz to 218 Hz [16]. With \( \Delta_2 = \Delta_e = 0 \), loss is suppressed at the center of the broad peak [see Fig. 3(a)]. For \( \Delta_2 = \Delta_1 + 2\mu_B(543.2 - B'_{\text{res}}) \), the loss is suppressed at the center of the shifted narrow peak; see Figs. 3(b) and 3(c). The continuum-dressed-state model (the solid red curves) predicts the features for all three data sets using the same Rabi frequencies, \( \Omega_1 \) and \( \Omega_2 \), which are close to the predicted values [16]. We note that the predicted central peaks in Fig. 3(c) are somewhat larger than the measured values, which may arise from jitter in the two-photon detuning and the intensity variation of the \( \omega_2 \) beam across the atom cloud.

We have also measured light-induced loss and loss suppression as a function of the \( \omega_1 \) laser frequency near the broad resonance. Data for \( B = 840 \) G G are shown in Fig. 4, where \( \Delta_e = 10 \) MHz. Since the maximum suppression occurs for \( \Delta_e = \Delta_2 \), we observe an asymmetric loss suppression window. By measuring the number of atoms as a function of time with the magnetic field tuned to the suppression point (see Fig. 5), we observe an increase of the inelastic lifetime near the broad resonance of \( ^{7}\text{Li} \) from 0.5 ms with a single laser field to 400 ms with the two-field method, despite the large background scattering length of \(-1405 \delta_B [22,23]\) limited by jitter in the two-photon detuning.

Measured light-induced loss spectra are compared to predictions by calculating the atom loss rate [16]. For a 50–50 mixture of two hyperfine states, the total density decays according to \( \dot{n}(r,t) = -\frac{1}{2} \langle K_2(k,r) \rangle \langle n(r,t) \rangle \). Here, the angle brackets in \( \langle K_2(k,r) \rangle \) denote an average over the relative-momentum \( \hbar k \) distribution. As the Rabi frequencies that determine \( K_2 \) generally vary in space, we include an additional position dependence in \( \langle K_2 \rangle \). For simplicity, we assume in this Letter a classical Boltzmann distribution, which is applicable in the high temperature regime employed in the measurements and defer treatment of quantum degeneracy and many-body effects to future work.

The loss rate constant \( K_2(k) \) is calculated from the optically modified scattering state in the continuum-

FIG. 3. Loss suppression using the two-field optical technique. The \( \omega_1 \) beam shifts the narrow Feshbach resonance. The frequency of the \( \omega_2 \) beam is chosen to suppress loss from (a) the broad resonance and (b) the shifted narrow Feshbach resonance. (c) Expanded view of (b) near the loss suppression region. Pulse duration \( \tau = 5.0 \) ms; \( T = 4.5 \) \( \mu \)K; \( \Omega_1 = 2.6\gamma_e \); \( \Omega_2 = 0.8\gamma_e \); \( \Delta_2 = \Delta_1 = 30.2 \) MHz; \( B'_{\text{res}} = 541.1 \) G. Vertical dashed line, position of the unshifted narrow resonance. Solid red curves, continuum-dressed-state model [16].

FIG. 4. Loss suppression near the broad resonance at 832.2 G, for \( B = 840 \) G, as a function of single photon detuning by sweeping the \( \omega_1 \) laser frequency. Pulse duration \( \tau = 5.0 \) ms; \( T = 14.8 \) \( \mu \)K; \( \Omega_1 = 1.36\gamma_e \); \( \Omega_2 = 0.9\gamma_e \); \( \Delta_2 = 10.0 \) MHz. Maximum suppression occurs for \( \Delta_e = \Delta_2 = 10.0 \) MHz, where \( \delta_e = 0 \). Solid red curve, continuum-dressed-state model [16].
interaction of the colliding atom pair with the optical fields is described in the bare-state basis [see Fig. 6(a)], with the singlet states $|g_1\rangle$, $|g_2\rangle$, and $|e\rangle$ and the triplet continuum $|T,k\rangle$. Using the continuum-dressed-state basis [see Fig. 6(b)], the bare states $|g_1\rangle$ and $|T,k\rangle$ are replaced by the dressed bound state $|E\rangle$ and the Feshbach resonance scattering state $|E_k\rangle$. These dressed states already contain the hyperfine coupling constant $V_{HF}$, permitting consistent adiabatic elimination of the excited state amplitude $|e\rangle$, even for broad Feshbach resonances where $V_{HF}$ is large compared to $\gamma_r$. From the scattering state, we determine the corresponding two-body scattering amplitude $f(k)$, which yields $K_2(k)$ from the inelastic cross section. The new model shows that the light shifts arising from the $\Omega_2$ beam have a different relative-momentum dependence for broad resonances than for narrow resonances and reproduce previous calculations [7,12,13] that are valid only for narrow resonances. For broad resonances, the new model avoids a fixed loss point at $B_{\sigma_0}$, which is incorrectly predicted by narrow resonance models [7,12,13].

In conclusion, we have demonstrated that closed-channel dark-state methods can produce large shifts of magnetic Feshbach resonances and strong suppression of spontaneous scattering, enabling flexible control of the trade-off between loss and tunability. We have established a continuum-dressed-state model that fits our measured loss spectra in shape and magnitude, using the measured Rabi frequencies and trap parameters. A key result of this model is the prediction of the relative-momentum dependence of the light-induced level shifts for both broad and narrow resonances, resolving a long-standing issue with predictions for broad resonances. Using the predicted relative-momentum dependence of the scattering amplitude, the model predicts optical control of not only the scattering length [16] but also the effective range [16], which will be experimentally studied in future work.

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